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STUDIES OF LASER INDUCED TURBULENCE IN
THE ATMOSPHERE III

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FINAL REPORT

STUDIES OF LASER INDUCED

TURBULENCE IN THE ATMOSPHERE III*

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SUMMARY OF RESEARCH

Various aspects of the propagation of an intense laser beam through the atmosphere are considered. The basic laser-fluid equations are presented and a linearized analysis of these equations is given which predicts a very low power threshold for Brueckner-Jorna-type convective instabilities. Another class of instabilities is predicted to be of more practical importance than the convective instabilities and an effective Reynolds number is derived which may help to characterize these turbulent instabilities. A computer solution of the full set of non-linear equations is described, and the concept of "utility analysis" of numerical differencing schemes is introduced. With the computation scheme used, the laser pulse could be followed for only 10^{-5} seconds; so enormous energy was put into the pulse to enhance the interaction with the fluid. Thus the initial pulse distortion could be observed. Analytical evaluation of the computer results produces a detailed quantitative check and suggests that a combination of analytic and numerical methods would allow a pulse to be conveniently followed for much longer periods of time.

INDEX OF TECHNICAL REPORTS

- W. G. Wagner, J. H. Marburger, R. F. Tooper, and B. J. Lambrigtsen
Studies of Laser Induced Turbulence in the Atmosphere I.
Report to Advanced Research Projects Agency, 1971.
- W. G. Wagner, J. H. Marburger, B. J. Lambrigtsen, and R. F. Tooper
Studies of Laser Induced Turbulence in the Atmosphere II.
Report to Advanced Research Projects Agency, 1972.
- W. G. Wagner, and J. H. Marburger Studies of Laser Induced
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Research Projects Agency, 1972.
- John D. Reichert Laser Beams in Flooming Wind Prisms. Report to
Advanced Research Projects Agency, June 1971.
- John D. Reichert, W. G. Wagner, and W. Y. Chen Propagation of
Intense Laser Beams in Air. Report to Advanced Research
Projects Agency, 1971
- Wei-Yeu Chen, John D. Reichert, and W. G. Wagner Utility
Analysis for Numerical Solution of Differential Equations.
Report to Advanced Research Projects Agency, 1971.
- Annual Summary Report dated October 1969 published as:
Computer Studies of High Intensity Beam Propagation
published in the Proceedings of the Laser Propagation
Meeting at Lincoln Laboratory, 15 and 16 October, 1969.

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J. D. Reichert, W. G. Wagner, and W. Y. Chen Instabilities of Intense Laser beams in Air J. Applied Physics, 41, 3641-46 (1973)

J. D. Reichert, W. G. Wagner, and W. Y. Chen Propagation of an Intense Gaussian Laser Pulse in Air J. Applied Physics, 41, 3647-58 (1973)

Instabilities of intense laser beams in air*

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Various aspects of the propagation of an intense laser beam through the atmosphere are considered. The basic laser-fluid equations are presented and a linearized analysis of these equations is given which predicts a very low power threshold for Brueckner-Jorna-type convective instabilities. Another class of instabilities is predicted to be of more practical importance than the convective instabilities and an effective Reynolds number is derived which may help to characterize these turbulent instabilities. In the following paper by the same authors a computer solution of the full set of nonlinear equations is described and the initial development of the laser-fluid interaction is investigated.

I. INTRODUCTION

The present paper and the following paper¹ are concerned with the distortions of a laser beam produced by density and thermal variations in a fluid medium. These distortions have been the subject of many investigations which can be classified into two groups, depending upon whether time dependence is considered. Most available experiments are conveniently understood by reference to theoretical studies of the gross effects of thermal deposition and fluid motion which assume that a steady state will be achieved for the deflection and distortion of the laser beam. On the other hand Brueckner and Jorna² have discovered that some of the solutions for beam propagation are unstable so that under certain conditions a steady state may not develop. The Brueckner-Jorna instabilities were discovered in a linearized analysis, but the threshold for such instabilities has not been discussed previously. The Brueckner-Jorna instabilities are not observed in practice because they are convective instabilities and cannot develop within typical distances allowed for the propagation of beams in the laboratory.

In the present paper the linearized analysis is presented, keeping the four-photon coupling induced by periodic fluctuations in the dielectric constant, so that among other things the threshold behavior will be exposed. The basic equations are given in Sec. II and the linearized analysis is presented in Sec. III. In Sec. IV another class of instabilities is predicted to be of more practical importance than the convective instabilities and an effective Reynolds number is derived which may help to characterize these turbulent instabilities.

The companion paper¹ describes a computer solution of the full set of nonlinear laser-fluid equations. Enormous power was presumed for the laser beam in order to drive the laser-fluid interaction as fast as possible in an attempt to watch the onset of distortions of the beam.

II. BASIC EQUATIONS

When an intense laser beam propagates through a fluid, many interesting phenomena take place. This laser-fluid system can be described by a macroscopic model which involves Maxwell's equations, the Navier-Stokes equation, an energy conservation equation, and the continuity equation for fluid motion. These equations, which de-

scribe the behavior of intense electromagnetic beams and the associated sound and thermal fluctuations, are coupled by stimulated Raman scattering, electrostriction, the high frequency Kerr effect, absorption heating, and the density and temperature dependence of the dielectric constant. In this paper a systematic discussion is presented for an intense laser beam propagating through air, which has a negligible Kerr constant. If the frequencies are outside the Raman scattering range, the instabilities are primarily caused by optical-acoustic coupling of the laser beam and the gases. These effects are of long duration compared to those of self-focusing. As the beam passes through air, the intensity profile induces a nonuniform temperature gradient transverse to the propagating direction of the beam, due to the energy absorption from the beam. This thermal nonequilibrium and electrostriction together cause the generation of a density gradient and hence a sound wave. These density changes react back on the incident beam through changes in the dielectric constant.

The equations describing propagation of electromagnetic radiation and the equations describing fluid behavior are widely known.³⁻⁵ Dropping unimportant terms from the full equations, we take the following set of nonlinear coupled partial differential equations for description of the macroscopic representation of the laser-fluid system:

Wave equation and Clausius-Mosotti relation:

$$\nabla^2 \mathbf{E} = \frac{1}{c^2} \frac{\partial^2}{\partial t^2} (\epsilon \mathbf{E}) + \frac{\alpha}{c} \frac{\partial}{\partial t} (\sqrt{\epsilon} \mathbf{E}), \quad (1)$$

$$\epsilon = \epsilon_L + \epsilon_2 \langle E^2 \rangle_{av},$$

$$\rho \left(\frac{\partial \epsilon}{\partial \rho} \right)_T = \frac{(\epsilon - 1)(\epsilon + 2)}{3}, \quad \left(\frac{\partial \epsilon}{\partial T} \right) \neq 0; \quad (2)$$

Navier-Stokes equation:

$$\rho \frac{D\mathbf{v}}{Dt} = \rho \mathbf{g} - \nabla P + \mathbf{f}_{es} + \mathbf{f}_{visc}, \quad (3)$$

$$\mathbf{f}_{visc} = \eta \nabla^2 \mathbf{v} + (\eta + \eta') \nabla (\nabla \cdot \mathbf{v}), \quad (4)$$

$$\mathbf{f}_{es} = \frac{1}{2} \nabla \left\langle \left[E^2 \rho \left(\frac{\partial \epsilon}{\partial \rho} \right)_T \right] \right\rangle_{av} - \frac{1}{2} \langle E^2 \nabla \epsilon \rangle_{av}; \quad (5)$$

heat transfer equation:

$$\rho C_v \frac{DT}{Dt} - \frac{C_v(\gamma-1)}{\beta} \frac{D\rho}{Dt} = \phi_\eta + \nabla \cdot (\kappa \nabla T) + \alpha c \sqrt{\epsilon} \langle E^2 \rangle_{av}, \quad (6)$$

$$\phi_\eta = \left[\eta \left(\frac{\partial v_i}{\partial x_j} + \frac{\partial v_j}{\partial x_i} \right) + \eta' \frac{\partial v_i}{\partial x_i} \delta_{ij} \right] \frac{\partial v_i}{\partial x_j}; \quad (7)$$

fluid continuity equation:

$$\frac{\partial \rho}{\partial t} + \nabla \cdot (\rho \mathbf{v}) = 0; \quad (8)$$

equation of state:

$$P = P(\rho, T) = R\rho T. \quad (9)$$

in the wave equation the term involving α , the linear absorption coefficient, is associated with a model for the absorption of electromagnetic energy by the fluid.⁶ The absorption coefficient is taken to be independent of the frequency of the electromagnetic radiation, so the model is not valid near the resonance lines of the molecules in the fluid. Also the model does not include kinetic rate equations, so saturation effects are not considered. In the present model the electric field \mathbf{E} will be damped by a factor $\exp(-\frac{1}{2}\alpha z)$, where z is the direction of propagation and the energy deposited in the medium is taken to be $\alpha c \sqrt{\epsilon} \langle E^2 \rangle_{av} = \alpha I_L$, where I_L is the laser intensity in erg/sec cm².⁷

The terms ϵ_L and ϵ_2 are the linear and nonlinear permittivity coefficients, respectively, ρ is the fluid mass density, and \mathbf{v} is the velocity of a "material element" of the fluid. The convective derivative D/Dt follows the motion of a material "particle" of the fluid relative to a fixed coordinate system and is expressed in the form

$$\frac{D(\quad)}{Dt} = \frac{\partial(\quad)}{\partial t} + \mathbf{v} \cdot \nabla(\quad).$$

The vector \mathbf{g} is the gravitational acceleration vector and η and η' are the shear and compressional viscosity coefficients, respectively.

The electrostrictive force density \mathbf{f}_{ee} is given by

$$(\mathbf{f}_{ee})_i = \left\langle \frac{\partial}{\partial x_j} \sigma_{ij} \right\rangle_{av},$$

where σ_{ij} is the interaction stress tensor for the electromagnetic field and the fluid and the angular brackets indicate a time average over several optical periods. A derivation of the stress tensor is given on p. 67 of Ref. 4:

$$\sigma_{ij} = -\frac{1}{2} E^2 \left[\epsilon - \rho \left(\frac{\partial \epsilon}{\partial \rho} \right)_T \right] \delta_{ij} + \epsilon E_i E_j,$$

but this tensor is not strictly correct for optical fields because an isothermal constraint was imposed. A similar derivation with an isentropic constraint gives the same result, except that the partial derivative $(\partial \epsilon / \partial \rho)_T$ at constant temperature is replaced by $(\partial \epsilon / \partial \rho)_s$, the derivative at constant entropy. The difference in these two constraints is contained in the thermodynamic relation

$$\rho \left(\frac{\partial \epsilon}{\partial \rho} \right)_s = \rho \left(\frac{\partial \epsilon}{\partial \rho} \right)_T + T \left(\frac{\partial \epsilon}{\partial T} \right)_\rho \left(\frac{\gamma C_p}{\beta v_s^2} - \beta T \right)^{-1}. \quad (10)$$

The difference term in Eq. (10) is very small because,

for gases, $(\partial \epsilon / \partial T)_\rho \approx 0$. The term $(\gamma C_p / \beta v_s^2 - \beta T)^{-1}$ contributes a factor of roughly $\frac{1}{2}$. Actually, neither constraint is strictly valid, but corrections would be small and would necessitate a detailed examination of fluid boundary layers and the explicit mechanisms of heat deposition in the control volume.

The thermodynamic quantities appearing in the above equations are C_v and C_p , the specific heats in erg/g deg at constant volume and pressure, respectively, and γ is the ratio of specific heats, C_p / C_v ; β is the thermal expansion coefficient, $-(1/\rho)(\partial \rho / \partial T)_p$; v_s is the isentropic velocity of sound, $[(\partial P / \partial \rho)_s]^{1/2}$; and κ is the thermal conductivity of the fluid.

Due to the complexity of the laser-fluid equations shown above, it is not possible to obtain exact solutions analytically. The linearized solutions have been discussed² and a number of computer solutions have recently been given by various groups.⁸ In Sec. III a linearized analysis of this set of equations is presented.

III. LINEARIZED ANALYSIS

Linearized analysis is a standard perturbation technique. In this scheme it is assumed that each of the dependent variables in the problem can be expressed as the sum of its slowly varying zeroth order component and a small first-order correction.⁹ In this way, a set of linear equations for small disturbances is obtained. This approach to the analysis of the laser-fluid system was first investigated by Brueckner and Jorna.² In the present approach, two variables are used to describe the perturbed electromagnetic field, one for the component of the field which is vibrating in phase with the primary beam and one for the component out of phase. In this way, the four-photon coupling induced by periodic fluctuations in the dielectric constant can be included. This coupling was not included in the original formulation given by Brueckner and Jorna and accounts for the absence of a threshold in their analysis. The dispersion relation for these linearized equations has been evaluated and is more complicated in structure than that presented by Brueckner and Jorna. For propagation through air, however, the numerical differences are minor. The wave with the largest growth rate, resulting from resonant interactions between scattered electromagnetic waves and the thermal wave, propagates almost perpendicularly to the laser beam. The direction is such that the change in frequency of the scattered electromagnetic wave and the frequency of the thermal wave (which is zero) are approximately the same.

Writing

$$\begin{aligned} \mathbf{E} &\approx \mathbf{E}_{(0)} + \mathbf{E}_{(1)}, \\ \rho &\approx \rho_0 + \rho_1, \\ T &\approx T_0 + T_1, \\ \epsilon_{0e} &\approx \epsilon_{L(0)} + \epsilon_{2(0)} \langle E_{(0)}^2 \rangle_{av}, \end{aligned} \quad (11)$$

and taking

$$\mathbf{E}_{(0)} = \frac{1}{2} \hat{\mathbf{e}}_y E_0 \exp[i(\omega_L t - k_L z)] \exp(-\frac{1}{2}\alpha z) + \text{c. c.}, \quad (12)$$

$$\begin{aligned} \mathbf{E}_{(1)} &= \frac{1}{2} \hat{\mathbf{e}}_y \{ f \exp[i(\omega_s t - \mathbf{k}_s \cdot \mathbf{x})] \\ &\quad + g \exp[-i(\omega_s t - \mathbf{k}_s \cdot \mathbf{x})] \} \exp(-\frac{1}{2}\alpha z) + \text{c. c.}, \end{aligned} \quad (13)$$

$$\left(\frac{\rho}{\Gamma}\right) = \frac{1}{2} \left(\frac{\rho'}{\Gamma'}\right) \exp[i(\omega t - \mathbf{k} \cdot \mathbf{x})] + \text{c. c.}, \quad (14)$$

where

ω_L = laser frequency,

$\omega_s = \omega_L \pm \omega$,

$\mathbf{k}_s = k_L \hat{\mathbf{e}}_L \pm \mathbf{k}$,

and treating the velocity \mathbf{v} to be a first-order quantity, we obtain from Eqs. (1)–(9) the following four equations relating f , g , ρ' , and T' :

$$\begin{aligned} &[(\epsilon_{0e} + \frac{1}{2}\epsilon_{2(0)}E_0^2)\omega_s^2 - c^2k_s^2]f + (\frac{1}{2}\epsilon_{2(0)}E_0^2\omega_s^2)g \\ &+ (\frac{1}{2}AE_0\omega_s^2)\rho' + (\frac{1}{2}BE_0\omega_s^2)T' = 0; \end{aligned} \quad (16)$$

$$\begin{aligned} &(\frac{1}{2}\epsilon_{2(0)}E_0^2\omega_s^2)f + [(\epsilon_{0e} + \frac{1}{2}\epsilon_{2(0)}E_0^2)\omega_s^2 - c^2k_s^2]g \\ &+ (\frac{1}{2}AE_0\omega_s^2)\rho' + (\frac{1}{2}BE_0\omega_s^2)T' = 0; \end{aligned} \quad (17)$$

$$\begin{aligned} &[\alpha c(\epsilon_{0e})^{1/2}E_0]f + [\alpha c(\epsilon_{0e})^{1/2}E_0]g + \{i[C_s(\gamma - 1)/\beta]\omega\}\rho' \\ &+ (-\kappa k^2 - i\rho_0 C_s \omega)T' = 0; \end{aligned} \quad (18)$$

$$\begin{aligned} &[\frac{1}{\eta}(\epsilon_{0e} - 1)(\epsilon_{0e} + 2)E_0 k^2]f + [\frac{1}{\eta}(\epsilon_{0e} - 1)(\epsilon_{0e} + 2)E_0 k^2]g \\ &+ \{\omega^2 - iN\omega k^2 - [u^2 - \frac{1}{\eta}(\epsilon_{0e} - 1)AE_0^2 k^2]\rho' \\ &+ \{[-u^2\beta\rho_0 + \frac{1}{\eta}(\epsilon_{0e} - 1)BE_0^2]k^2\}T' = 0. \end{aligned} \quad (19)$$

Throughout these equations, E_0 has been written in place of $|E_0|$. Since only the magnitude of the complex amplitude appears, E_0 may be considered to be real and positive without loss of generality. For convenience in writing these equations, the following notation has been introduced:

$$u^2 \equiv v_s^2/\gamma, \quad (20)$$

so that u is the isothermal speed of sound in the unperturbed medium, and

$$\begin{aligned} A &\equiv \left(\frac{\partial \epsilon_L}{\partial \rho}\right)_T + \left(\frac{\partial \epsilon_2}{\partial \rho}\right)_T \langle E_{(01)}^2 \rangle_{\text{av}} \approx \frac{1}{\rho_0} \frac{(\epsilon_{0e} - 1)(\epsilon_{0e} + 2)}{3}, \\ B &\equiv \left(\frac{\partial \epsilon_L}{\partial T}\right)_\rho + \left(\frac{\partial \epsilon_2}{\partial T}\right)_\rho \langle E_{(01)}^2 \rangle_{\text{av}}, \\ N &\equiv (2\eta + \eta')/\rho_0. \end{aligned} \quad (21)$$

The consistency condition for the linearized equations is the vanishing of the determinant of the coefficients of f , g , ρ' , and T' . Thus we obtain the following relation between the frequency ω and the wave vector \mathbf{k} , the dispersion relation for the system:

$$\begin{aligned} &\{(\omega - i\kappa'k^2)[\omega^2 - iN\omega k^2 - (u^2 - A')k^2] - (\gamma - 1)(u^2 - B')\omega k^2\} \\ &\times [\xi_s \xi_s - \frac{1}{2}\epsilon_{2(0)}E_0^2(\xi_s \omega_s^2 + \xi_s \omega_s^2)/2\omega_L] \\ &= -[Ak^2[\frac{1}{2}\rho_0 A(\omega - i\kappa'k^2) + i(\alpha c n_{0e} \beta/C_s)(u^2 - B')]] \\ &+ B\{[(\gamma - 1)/\beta]\frac{1}{2}AE_0 k^2 + i(\alpha c n_{0e}/\rho_0 C_s)[\omega^2 - iN\omega k^2 \\ &- (u^2 - A')k^2]\}\frac{1}{2}E_0^2(\xi_s \omega_s^2 + \xi_s \omega_s^2)/2\omega_L. \end{aligned} \quad (22)$$

In order to put Eq. (22) in the slightly more compact form shown above, the following additional abbreviations have been introduced:

$$\kappa' \equiv \kappa/\rho_0 C_s,$$

$$\eta_{0e} \equiv (\epsilon_{0e})^{1/2} = \text{index of refraction of unperturbed medium,}$$

$I_L \equiv \frac{1}{2}\eta_{0e}cE_0^2$ = power in incident laser beam/unit area,

$$A' \equiv \frac{1}{\eta}(\epsilon_{0e} - 1)AE_0^2, \quad (23)$$

$$B' \equiv \frac{1}{\eta}(\epsilon_{0e} - 1)BE_0^2/\rho_0\beta,$$

$$2\omega_L \xi_s \equiv c^2 k_s^2 - \epsilon_{0e} \omega_s^2,$$

$$2\omega_L \xi_s \equiv c^2 k_s^2 - \epsilon_{0e} \omega_s^2.$$

For comparison, the dispersion relation obtained by Brueckner and Jorna² for frequencies outside the Raman scattering range is

$$\omega(\omega^2 - v_s^2 k^2 - iN\omega k^2)\xi_L + I_L(\nu\omega + i\delta)k^2 = 0, \quad (24)$$

where

$$\begin{aligned} \xi_L &\equiv [k^2 - (2\epsilon_{2(0)}/n_{0e}^3 c)I_L k_L^2] - 4(n_{0e} k_L/c k)[\omega - (c/n_{0e})k_s]^2, \\ \nu &\equiv \rho_0 A^2 k_L^2/n_{0e}^3 c, \end{aligned} \quad (25)$$

$$\delta \equiv 2A\beta n_{0e}^3 \alpha k_L^2/n_{0e}^3 C_s.$$

Although Eq. (22) is considerably more complicated in structure than Eq. (24), the general features of the two equations are the same. As a first approach to the analysis of (22), one should realize that the power in the primary laser beam is proportional to E_0^2 . Thus, the free modes of the system can be obtained by letting $E_0^2 \rightarrow 0$. With no power in the incident beam, therefore, (22) reduces to

$$\begin{aligned} &\{[\omega - i(\kappa/\rho_0 C_s)k^2][\omega^2 - iN\omega k^2 - u^2 k^2] - (\gamma - 1)u^2 \omega k^2\} \\ &\times [c^2(k_L^2 + k^2 + 2k_L k_s) - \epsilon_{L(0)}(\omega_L + \omega)^2] \\ &\times [c^2(k_L^2 + k^2 - 2k_L k_s) - \epsilon_{L(0)}(\omega_L - \omega)^2] = 0. \end{aligned} \quad (26)$$

The first factor in (26) contains a nonpropagating thermal wave and two damped sound waves coupled by the term $(\gamma - 1)u^2 \omega k^2$. The last two factors correspond to the four free modes for scattered undamped electromagnetic waves:

$$\omega/\omega_L = \sigma \pm [1 + k^2/k_L^2 + 2\sigma k_s/k_L]^{1/2}, \quad (27)$$

where $\sigma = \pm 1$, specified by given values of k^2 and k_s . Two of the roots are low frequency ($\omega \ll \omega_L$), whereas the other two have frequencies of the same order as the laser frequency. It is clear that the roots at the high frequency should be eliminated, because it has been assumed previously, in evaluating time averages, that the perturbed solutions vary much more slowly than the optical waves. Therefore, factors like

$$c^2 k^2 - \epsilon_{L(0)} \omega_s^2 \quad (28)$$

will be replaced by

$$\epsilon_{L(0)}(2\omega_L)(c_L k_s - \omega_s), \quad (29)$$

where

$$c_L \equiv c/[\epsilon_{L(0)}]^{1/2} = \text{velocity of light in the medium.}$$

Thus the free modes will include three thermal-sound waves and two electromagnetic waves.

All of the terms in (26) result from the left-hand side of (22) because the right-hand side is proportional to the power in the primary laser beam. Now, as the power in the laser beam is turned on, the right-hand side couples the five free modes described by (26). Additional tiny coupling arises inside the left-hand side itself through the A' , B' , and $\epsilon_{2(0)}$ terms.

For detailed consideration, the case of a primary laser beam at 10.6 μ propagating through air at approximate-ly 10°C and at standard pressure will be discussed. The numerical values for the parameters appearing in the dispersion relation are¹⁰

$$\begin{aligned}\omega_L &= 1.773 \times 10^{14} \text{ sec}^{-1}, \\ k_L &= 5920 \text{ cm}^{-1}, \\ \rho_0 &= 1.25 \times 10^{-3} \text{ g/cm}^3, \\ N &= 0.284 \text{ cm}^2/\text{sec}, \\ \beta &= 3.67 \times 10^{-3} \text{ deg}^{-1}, \\ C_v &= 7.143 \times 10^6 \text{ erg/g deg}, \\ \kappa' &= 0.28 \text{ cm}^2/\text{sec}, \\ n_0 &= 1 + 2.82 \times 10^{-4}, \\ u^2 &= 8.39 \times 10^8 \text{ (cm/sec)}^2, \\ \gamma &= 1.4, \\ \epsilon_{L(0)} &= 1 + 5.65 \times 10^{-4}, \\ \alpha &= (3 \times 10^{-7} \text{ cm}^{-1}) \alpha_0, \\ \alpha \eta_0 \epsilon u^2 \beta / C_v &= 3.831 \times 10^3 \alpha_0, \\ \alpha \eta_0 \epsilon / \rho_0 C_v &= 0.988 \alpha_0, \\ \frac{1}{2} \epsilon_{2(0)} E_0^2 &= (1.37 \times 10^{-25} \text{ sec}^3/\text{g cm}^2) I_L, \\ A &= 0.452 \text{ cm}^3/\text{g} + (1.10 \times 10^{-22} \text{ cm sec}^3/\text{g}^2) I_L, \\ A(\gamma - 1)/2\beta &= 24.7 \text{ cm}^3 \text{ deg/g} \\ &\quad + (6.00 \times 10^{-21} \text{ cm sec}^3 \text{ deg}^{-1} \text{ g}^2) I_L, \\ \epsilon_{0e} &= 1 + 5.65 \times 10^{-4} + (1.37 \times 10^{-25} \text{ sec}^3/\text{g cm}^2) I_L, \\ A' &= (2.84 \times 10^{-15} \text{ cm}^2 \text{ sec/g}) I_L + (1.38 \times 10^{-36} \text{ sec}^4/\text{g}^2) I_L^2, \\ B &= \epsilon_{2(0)} E_0^2 / T_0 = -(4.84 \times 10^{-26} \text{ sec}^3/\text{g cm}^2 \text{ deg}) I_L, \\ B' &= -(0.662 \times 10^{-36} \text{ sec}^4/\text{g}) I_L^2.\end{aligned}\quad (30)$$

In the above list a dimensionless absorption constant α_0 of order unity has been introduced and the power I_L is in units of erg/sec per cm^2 . Now, using these numerical values, one finds that the power-dependent terms are very small for power fluxes less than 10 MW/ cm^2 , except for the term which represents energy absorption. In other places in the dispersion relation, the power-dependent terms are connected with the nonlinear index, and will be omitted in the following. (The terms omitted are related to self-focusing in a manner described by Brueckner and Jorna.²) This neglect of the nonlinear index and of the weak dependence of the optical coefficients of gases on the temperature for fixed density allows the simplification of the dispersion relation to

$$\begin{aligned}[(\omega - i\kappa'k^2)(\omega^2 - iN\omega k^2 - u^2k^2) - (\gamma - 1)u^2\omega k^2] \\ \times (c_L k_- - \omega_+)(c_L k_- - \omega_-) \\ = -(Ak^2\omega_L I_L / 2\alpha_L)[(c_L k_- - \omega_+) + (c_L k_- - \omega_-)] \\ \times \left\{ \frac{1}{2} \rho_0 A(\omega - i\kappa'k^2) + i\alpha \eta_0 \epsilon u^2 \beta / C_v \right\}.\end{aligned}\quad (31)$$

In addition, for air at reasonable powers, the term $\frac{1}{2} \rho_0 A(\omega - i\kappa'k^2)$ on the right-hand side is negligible compared to $\alpha \eta_0 \epsilon u^2 \beta / C_v$. Introducing the variable

$$\nu \equiv k_3/k, \quad (32)$$

instead of k_3 , and defining

$$\tau \equiv A\omega_L \beta \alpha I_L / \sqrt{\epsilon_L} C_v, \quad (33)$$

which corresponds to the power parameter used in Ref. 2, the dispersion relation can be written in the form

$$\begin{aligned}(\omega - i\kappa'k^2)(\omega^2 - iN\omega k^2 - u^2k^2) - (\gamma - 1)u^2\omega k^2 \\ = \frac{1}{2} i \tau \nu^2 k^2 \{ [\omega - \frac{1}{2} i \alpha c_L - c_L(k\nu + k^2/2k_L)]^{-1} \\ - [\omega - \frac{1}{2} i \alpha c_L - c_L(k\nu - k^2/2k_L)]^{-1} \}.\end{aligned}\quad (34)$$

The problem at this stage is the determination of the maximum growth rate of any Fourier component of a distortion of the plane wave as a function of the absorbed power from the beam. That is, one must solve the dispersion relation for the frequencies as a function of k , ν , I_L , and the characteristic parameters of the medium, and then find the maximum value of $-\text{Im}\omega$ for real ν and k with $|\nu| < 1$. Such a problem cannot be solved analytically without further approximations. One region of interest would be the high power limit, where the driving term would overwhelm the losses resulting from thermal conduction, viscosity, and the absorption of electromagnetic energy. In that case, all the imaginary terms in Eq. (34), with the sole exception of the i immediately preceding τ , can be dropped, reducing the dispersion relation to the form employed by Brueckner and Jorna in Eq. (45) of Ref. 2.

To proceed analytically, Brueckner and Jorna neglected the term in the second set of brackets on the right-hand side of (34), and assumed that the maximum growth rate would occur somewhere on the curve in the ν, k plane determined by the constraint

$$\text{Re}[\omega - c_L(k\nu + k^2/2k_L)] = 0. \quad (35)$$

Along that curve, the maximum growth rate is

$$(-\text{Im}\omega)_{\text{max}} = \frac{1}{2} \sqrt{\tau} (1.08), \quad (36)$$

which corresponds to

$$\nu, k = \left(\frac{1}{2} \tau \right)^{1/2} \left(\frac{0.97515}{0.56305} \right) \left(\frac{1}{0.93063} \right)^{1/2}. \quad (37)$$

(These results differ from those in Ref. 2, which are erroneous.) There is no assurance that the actual maximum growth rate does lie along the one-dimensional subset of the ν, k plane assumed in Ref. 2. We have conducted a search along the line

$$\nu + k/2k_L = 0.$$

However, the result for the maximum,

$$\frac{1}{2} \sqrt{\tau} (1.06), \quad (38)$$

is 2% smaller. No other curve in the ν, k plane has been found which allows an analytical search. Nevertheless, one suspects that these answers are sufficiently close and that further analytical effort is not justified, because of the previous approximations.

An interesting unknown not discussed previously is the power flux required to stimulate these instabilities. This threshold power is clearly a critical function of the losses in the system, which therefore renders it important to treat them carefully. If the second term on the right-hand side of Eq. (34) is dropped, the instabil-

ity appears to have no threshold, because the conduction loss, which must be overcome, vanishes as $k \rightarrow 0$. However, as $k \rightarrow 0$, the Stokes and anti-Stokes terms on the right-hand side of Eq. (34) tend to cancel each other, and, therefore, there is a threshold power flux for these stimulated thermal Rayleigh scattering instabilities. Using (34), a computer search for this threshold was performed and led to the result.

$$(I_L)_{\text{threshold}} = 0.329 \text{ mW/cm}^2. \quad (39)$$

This threshold was located at $k = 0.04 \text{ cm}^{-1}$ with $\nu = 1.1 \times 10^{-6}$. The degeneracy in the value of ν occurs because of the symmetry property contained in (34), $\nu \rightarrow -\nu$ implies $i\omega \rightarrow (i\omega)^*$. The dispersion relation of Brueckner and Jorna contained no threshold for the convective instabilities and, therefore, suggested that instabilities might be present for extremely low beam intensities. The $\frac{1}{2}$ -mW per cm^2 threshold obtained in the present analysis is certainly small in relation to intensities available for experiments.

The presence of a wind does not alter the growth rates for distortions. This can be easily seen by considering the problem from a frame of reference moving with the fluid. A uniform beam remains a uniform beam in the moving frame, although its direction of propagation is shifted. This shift in direction has no effect upon the stability discussion. The very low convective instability threshold does not mean that such effects are easy to observe or are of practical importance. An extremely long, and carefully protected, optical path length would be required in order to see these instabilities grow to observable sizes. Indeed, such instabilities have never been observed in the laser-fluid system.

There is, however, another class of instabilities involving the balance between inertial terms and buoyancy forces in the fluid. Such instabilities are discussed in Sec. IV.

IV. BEAM-FED TURBULENCE

Many experiments have been reported for which theories assuming steady-state beam profiles, after initial transients die out, provide rather good explanations of the principal features. However, that is probably true only because these experiments are conducted at relatively low power fluxes. Theoretically, one expects a time-dependent state of the system because of the instabilities discussed earlier. Such instabilities are not observed in practice because they cannot develop within the distances allowed for the propagation of beams. However, alternate considerations for a beam of finite cross section suggest that the beam may drive the fluid into a time-dependent, or turbulent, state at powers which are not completely unreasonable.

It may be impossible to prove analytically that such a turbulent state develops, because the investigation of hydrodynamic stability is very difficult even for the simplest flows. However, an argument can be made from dimensional considerations, an approach that promises to be very useful. Namely, for the problem of a beam of radius a and power flux I passing through air, it is possible to estimate a parameter W , which plays the role of an effective Reynolds number for our

problem. It will be shown that the parameter W takes on values of the order of 30 000 for a beam with intensity $I = 1 \text{ kW/cm}^2$ of radius 1 m. Since it is known that some flows with Reynolds numbers substantially lower than 30 000 are turbulent, the flows for the laser-heated atmospheric path should also be expected to show significant time dependence.

Consider the equations of motion for the air and the equation governing heat transfer, which take the following form if it is assumed that the air can be assumed incompressible (that amounts to dropping terms of order u^2/v_s^2 , where u is a typical flow speed, and v_s is the speed of sound; for the problems under consideration, u^2/v_s^2 will be less than 10^{-5} , and the incompressible fluid approximation will be quite good):

$$\begin{aligned} \rho(\mathbf{v} \cdot \nabla)\mathbf{v} &= -\nabla P - \beta T_1 \rho \mathbf{g} + (\rho/\rho_0)\eta \nabla^2 \mathbf{v} \\ &\quad + (\rho/\rho_0)(\eta + \eta') \nabla(\nabla \cdot \mathbf{v}), \\ \rho(\nabla \cdot \mathbf{v}) &= -(\mathbf{v} \cdot \nabla)\rho = \beta \rho(\mathbf{v} \cdot \nabla)T_1, \\ \rho C_p(\mathbf{v} \cdot \nabla)T_1 &= \alpha I + \nabla \cdot (\kappa \nabla T_1) \\ &\quad + (\rho/\rho_0)[\eta(v_{1,1} + v_{1,1})^2 + \eta'(\nabla \cdot \mathbf{v})^2]. \end{aligned} \quad (40)$$

Assuming the Reynolds number is high, the inertial terms will dominate the viscous terms in the Navier-Stokes equation. Thus, there must be a balance between the inertial terms and the buoyancy forces, which implies that $\rho u^2/a \sim \alpha \beta T_1 g$, where ρ is the density of air, β is the coefficient of thermal expansion, T_1 is a typical value for the temperature rise, and g is the acceleration due to gravity. The pressure variation will be of order ρu^2 . In the heat-transfer equation, the convection term will dominate the conduction term, and the beam heating will overwhelm the viscous dissipation, so that there must be a balance between heat deposition from the beam and convective heat transfer, which implies that $\rho C_p u T_1/a \sim \alpha I$. Combining these two relations we find that $u^3 \sim \alpha \beta a^2 I g / \rho C_p$. With this expression for u , we then define a parameter W , which is expected to indicate regimes where steady flow and where time-dependent flow may be anticipated. W is an estimate of the relative importance of inertial terms to viscous terms in controlling the flow:

$$\begin{aligned} W &= a u \rho_0 / \eta \\ &= a(\rho_0/\eta)(\alpha \beta a^2 I g / \rho C_p)^{1/3}. \end{aligned} \quad (41)$$

For a beam with $I = 10^{10} \text{ erg/cm}^2 \text{ sec}$, $a = 100 \text{ cm}$, $T_0 = 273^\circ \text{K}$, and $\alpha = 3 \times 10^{-6} \text{ cm}^{-1}$, $W = 30\,000$.

For many experiments described in the literature, the values of W are much smaller, and thus one would not expect any turbulent fluid flow to be observed. For example, in the original experiment of Gordon, Leite, Moore, Porto, and Whinnery¹¹ the parameter W takes on a value about 10^{-3} , and in the more recent experiments of Smith and Gebhardt,¹² W is of order 10.

The parameter W introduced here is different from the Grashof number, which is referred to in some discussions of the convective flows set up by the absorption of energy from a laser beam.¹³ In fact, the conceptual basis for using the Grashof number in a discussion attempting to explain the transition between smooth flow and time-dependent flow seems less relevant because

the Grasshof number appears to be more sensible when the thermal buoyant forces are balanced by viscous forces. In the present discussions, the thermal buoyant forces are balanced by inertial effects. It turns out that the number W is essentially the square root of the Grasshof number.¹⁴

We are planning experiments to determine the critical value of W , W_{cr} , which determines the onset of turbulent convective flows for the geometry appropriate to laser beam transmission. It is also our aim to attempt a theoretical evaluation of this critical value. At the present time we can only speculate that W_{cr} may be between 10^3 and 10^4 .¹⁵ The theoretical approach appears fairly difficult because the question of the stability of flows even without heat sources has only been answered theoretically for very simple geometries.^{16,17} The question of stability for fluids which are heated or cooled appears to have been treated mainly for cases in which the fluid would be motionless, and has not been explored for a problem like the present one.^{17,18} The first part of that problem would be to determine a steady-state flow pattern for a fluid with a distributed heat source within a right circular cylinder with its axis aligned at some angle to the vertical. For the case of a horizontal cylinder of infinitely small radius, the flow pattern has been calculated by Yih.^{19,20}

Unfortunately, however, that solution is not of great value for the present problem because the size of the cylinder radius is a critical parameter. Nevertheless, it is expected that Yih's solution will assist in obtaining the asymptotic form of the steady-state flow at large distances from the laser beam cylinder. Once that time-independent flow pattern has been determined, the linearization of the hydrodynamic equations for perturbations from the flow pattern will lead to an eigenvalue problem, which eventually will yield a critical value for W . Ostrach²¹ suggests that the eigenvalue problem can be bypassed as the stability of fully developed natural convection flows can be found by using the appropriate velocity profile in the classical theory of hydrodynamic stability. This assertion rests upon his analysis of the stability of free convection above a flat heated plate, where instability first appears for a Reynolds number of 283.

Above the threshold for beam-induced turbulence, governed by W_{cr} , general arguments²² lead to a size for the smallest eddies, $a(W_{cr}/W)^{3/4}$. For a 1-m beam, if W_{cr} should be about 10^3 , then the eddies might have sizes as small as 7 cm for a power flux of 1 kW/cm². The associated density fluctuation would then be expected to result in considerably increased scattering of the beam.

The arguments presented here show that there are substantially more important sources of instability in the laser-fluid system than those discussed in earlier linearized analyses. It is felt that these fluid instabilities will be enhanced by their interaction with the scattering of the laser beam, because of the general result that instabilities in fluids result if the heating of the

fluid is greater in those regions where the density of the fluid is greater.²³

At the present time we can only outline the general nature of the effects to be expected above a critical power level. Much additional work clearly needs to be done, both of an experimental and theoretical nature.

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Propagation of an intense Gaussian laser pulse in air*

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A numerical solution of the nonlinear macroscopic laser-fluid equations for propagation of a Gaussian laser pulse in air is described. The concept of "utility analysis" of numerical differencing schemes is introduced. With the computation scheme used, the laser pulse could be followed for only 10^{-5} sec, so enormous energy was put into the pulse to enhance the interaction with the fluid. Thus the initial pulse distortion could be observed. Analytical evaluation of the computer results produces a detailed quantitative check and suggests that a combination of analytic and numerical methods would allow a pulse to be conveniently followed for much longer periods of time. The preceding paper by the same authors describes various types of instabilities to be anticipated for propagation over long periods of time or with large powers.

I. INTRODUCTION

The work described in this paper was motivated by the discovery of instabilities in the system of equations describing electromagnetic wave propagation and fluid dynamics. These instabilities are described in the preceding paper.¹ A numerical solution of the full set of macroscopic laser-fluid equations (without saturation effects) is presented for a Gaussian laser pulse of enormous energy. The computation scheme employed was not optimal and the pulse could be followed for only 10^{-5} sec. Consequently, enormous energy was put into the pulse to enhance the laser-fluid interaction and drive the onset of strong distortions and instabilities.

Other computer solutions of the laser-fluid equations have recently been given² by various groups for physically reasonable powers. Such calculations have typically dealt with steady-state beams without saturation effects, but have included the effects of gravity. More recently saturation effects have been added and pulse shape is being considered.

In the course of these studies it became apparent that there was some merit to introducing a new concept to judge the value of an algorithm for computing the solutions of a system of partial differential equations. This concept was called "utility", and will be discussed only briefly in Sec. II. The full details will be published elsewhere.³ The advantage of this concept is that it is relatively easy to apply to complicated systems of partial differential equations, whereas the stability concept leads to a very complicated procedure for deciding on the value of a numerical routine.

Speed and memory size in a computer place certain restrictions on one's ability to investigate phenomena in the laser-beam problem. In the attempt to calculate distortions of the type predicted by the linearized instability analysis, cylindrical symmetry was imposed on the problem in order to facilitate the computer calculation. Had this not been necessary, or had some other independent variable been eliminated rather than the angle about the beam axis, much more pronounced evidence of beam and fluid instabilities would likely have been observed for substantially lower powers, powers that may be achievable. Arguments supporting this proposition are contained in the preceding paper.¹

In Secs. III and V we present the results of a calculation of beam distortion for a very high intensity pulse propagating through air for several kilometers. Analytical arguments are advanced in Sec. IV which suggest that the qualitative features of the distortions are correct, which lends credence to the computer output. Then, using the computer results as a check, the analytical procedure is shown to be adequate for detailed quantitative calculations. By combining the analytical procedure developed with appropriate computer support, the pulse could be followed much longer than the 10^{-5} sec described here.

II. DESCRIPTION OF THE PROCEDURE

The laser-fluid equations were solved in the near-field region of a laser pulse, initially Gaussian in both r and z , propagating through air at 1 atm of pressure and at 10°C . A cylindrical geometry was used and cylindrical symmetry (no dependence on the angle ϕ) was preserved at the price of dropping the gravity term in the Navier-Stokes equation. Having cylindrical symmetry amounts to a considerable simplification in the problem, so that the inclusion of the free convection effects due to gravity was not attempted in this analysis. The problem described above amounts to a mixed initial-boundary problem. The initial configurations of the laser beam and the fluid are specified subject to certain boundary conditions at $r=0$ which must be satisfied at all times. Furthermore, the boundary condition at $z=0$ is time dependent, because the tail of the Gaussian must be fed into the spatial region. For the numerical solution, a spatial mesh of grid points or stations is used to represent the rz plane. At a given instant in time, the values of the various dependent variables are obtained at all of the stations. The difference equations are then employed with these values of the dependent variables to advance a step in time. This procedure is repeated over and over until the desired time interval has been traversed. An explicit difference scheme was used in this calculation, because such schemes are simplest to handle.

The two major difficulties in using numerical techniques to solve differential equations by computer are error growth and excessive computation time. In order to control the error growth, the utility criterion mentioned in Sec. I has been used. Furthermore, highly accurate

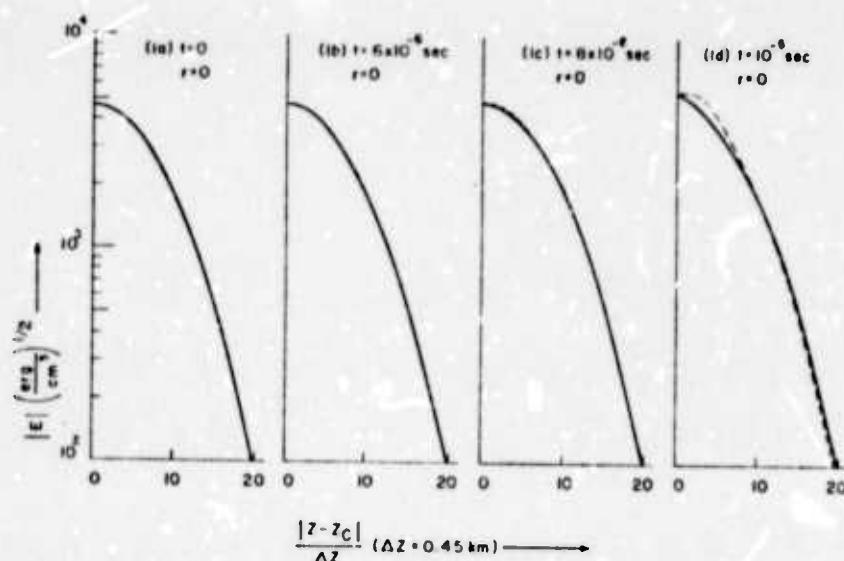


FIG. 1. On-axis magnitude of the electric field $|E|$ is shown as a function of $|z - z_c|$ at four different times. As shown in Eq. (16), $z_c \approx ct$ locates the "center" of the pulse. The solid curves show the leading edge of the pulse and the dotted curves depict the trailing edge. Distances along the z axis are expressed in units of the grid size; $\Delta z = 0.45$ km. The quantity $|E|$ shown is defined in Eq. (13), so that the exponential damping factor is not included in the graphs.

seven-point difference quotient representations of the differential operators were employed to reduce truncation error. In order to handle the economic problem of large computation time, a certain amount of efficiency is introduced by minimizing the amount of core storage required of the computer. This was accomplished in part by using overlaying techniques to store several pieces of information at the same site in the computer. Thus information is stored only as long as it is needed and then is replaced with current material. The computation time was also minimized by making use of a nonuniform grid. The seven-point difference relations allowed a relatively large grid size without undue truncation error and the nonuniform grid spacing permitted a greater grid density in the region of special interest. Thus an accurate solution could be obtained with a minimum of computation.

The laser-fluid equations are given in Eqs. (1)–(9) of Ref. 1. As mentioned above, the gravity term was dropped. Also, the thermal conductivity κ was taken to be constant because its derivatives are very small. The equation of state was taken to be the ideal gas law. The numerical values used for the various parameters are the same as those given for the linearized analysis in Ref. 1, because the same temperature and pressure were used for the undisturbed medium. The laser frequency ω_L and the dimensionless absorption constant α_0 were chosen to be $1.773 \times 10^{14} \text{ sec}^{-1}$ and 10, respectively. The wave equation for linearly polarized light in an absorptive medium is reduced to the scalar equation

$$c^2 \nabla^2 E = \frac{\partial^2}{\partial t^2} (\epsilon E) + \alpha \epsilon \frac{\partial}{\partial t} (\sqrt{\epsilon} E). \quad (1)$$

This equation is an approximate equation describing an electric field which is polarized linearly. Strictly speaking, of course, Maxwell's equations do not allow cylindrically symmetric linearly polarized beams in charge-free space.

The solution of (1) is taken to be in the form

$$E = \frac{1}{2} (E_1 + iE_2) \exp(i\omega_L t - k_L z) \exp(-\frac{1}{2}\alpha z) + \text{c. c.}, \quad (2)$$

where E_1 and E_2 are slowly varying functions of r and t and the laser frequency and wave number are related by

$$\epsilon_0 \omega_L^2 = c^2 k_L^2, \quad (3)$$

where ϵ_0 is defined in Eq. (11) of Ref. 1. Equation (2) was substituted into (1) and the second derivatives of E_1 and E_2 with respect to time were dropped. From the real and imaginary parts of the resulting equation, coupled simultaneous differential equations for E_1 and E_2 were obtained. These equations are also coupled to the differential equations for ρ , T , and the cylindrical components v_r and v_z of the fluid velocity. The use of cylindrical symmetry allows the simplification $v_\theta = 0$.

The easiest way to obtain a variable grid size is to introduce a transformation to a new independent variable. Thus, in order to have more grid points in the region of special interest, small r , the nonlinear transformation

$$r/r_0 = x/(1-x) \quad (4)$$

was employed. The scale value r_0 was chosen according to the dictates of convenience and will be shown in Eq. (9).

Because of the symmetry of the problem and the regularity of the differential equations, the following boundary conditions must be satisfied at $x=0$:

$$\begin{aligned} \frac{\partial E_1}{\partial x} = \frac{\partial E_2}{\partial x} = \frac{\partial T}{\partial x} = \frac{\partial \rho}{\partial x} = \frac{\partial v_r}{\partial x} = \frac{\partial v_z}{\partial x} = \frac{\partial^2 v_r}{\partial z^2} = v_r = 0 \\ \frac{\partial^2 v_r}{\partial x^2} = 2 \frac{\partial v_r}{\partial x}. \end{aligned} \quad (5)$$

The laser-fluid differential equations must, of course, be converted to difference equations before a computer solution can be attempted. As indicated earlier, seven-point difference quotients were used to represent the differential operators, but these difference quotients will not be presented. (The scheme was a straightforward explicit difference representation.)

Once the difference equations have been written, it is vital to have some criterion to determine useful time step sizes and corresponding grid spacings. It is not, however, necessary to demand that the difference scheme be stable in the classical sense. In fact, the classical notion of stability is irrelevant to computer solution of differential equations, particularly for non-

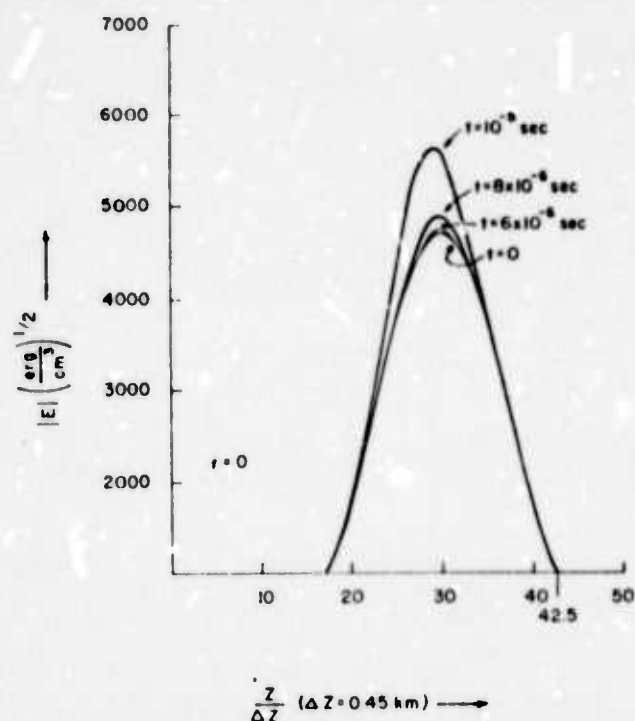


FIG. 2. For detailed comparison, the z profile of the on-axis electric field $|E|$ is shown at several times. The unit Δz is used for distances along the z axis. The curves have been displaced to the left and the leading edges made to coincide at height 1000 $(\text{erg}/\text{cm}^2)^{1/2}$ for $|E|$. The abscissa for this intersection of the curves has been labeled 42.5, the location of this point at $t = 0$.

linear systems. This argument is presented in detail in another paper³ and another kind of criterion, "utility", is proposed. A difference scheme is "useful" for prescribed ϵ and N if the computer solution of the scheme produces results for N time steps which differ in magnitude from the correct solution of the original differential equations by an amount less than ϵ . Thus, utility analysis is totally different in philosophy than stability and, incidentally, is far easier to apply. A number of theorems have been developed which allow rapid and powerful assessment of the utility regions of differencing schemes of all types. Because these matters would only extend the length of the present paper and because it is believed that utility procedures will be worthwhile in a broader context than the problem at hand, a separate paper is devoted to the subject and only a brief sketch of the analysis will be given here.

It turns out that a utility region for an explicit differencing scheme can be determined from a relation of the form

$$\Delta t < [(4\bar{\delta})(\frac{1}{100}N)]^{-1}, \quad (6)$$

where Δt is the time step size, N is the number of time steps to be made, $\epsilon \sim 10^{-3}$, and $\bar{\delta}$ is a function of the spatial mesh sizes and is determined from the differencing scheme. For the seven-point explicit difference scheme used to represent the laser-fluid system, $\bar{\delta}$ is trivially determined to be³

$$\bar{\delta} \approx c[3/k_L(r_0\Delta r)^2 + 2/\Delta z], \quad (7)$$

so that the utility restriction predicted is

$$1/c\Delta t \geq 12/k_L(r_0\Delta r)^2 + 8/\Delta z. \quad (8)$$

Since $\bar{\delta}$ is given only approximately by (7), the laser power I_L does not appear. The differencing scheme employed was not optimal and its utility region is entirely dominated by the electric field development and is insensitive to the fluid parameters. This was, however, the scheme employed in the present solution.

Condition (8) is such a strong constraint that one immediately wonders if it is really necessary to obey it. Part of the utility philosophy is to obtain a constraint like (8) and try it on the computer. Then one can try to violate the condition, using a larger Δt , and compare the results. This was done for the problem discussed in this paper and no escape from (8) was possible for the difference scheme used. In fact, if the criterion was violated by a factor of order 5 in Δt , then classic instability phenomena were observed in the computer output. Thus, by a stroke of bad luck, it appears that (8) must be obeyed.

In order to emphasize the implications of (8) for the study of the propagation of laser pulses, a description of the accessible parameter regime will now be given.

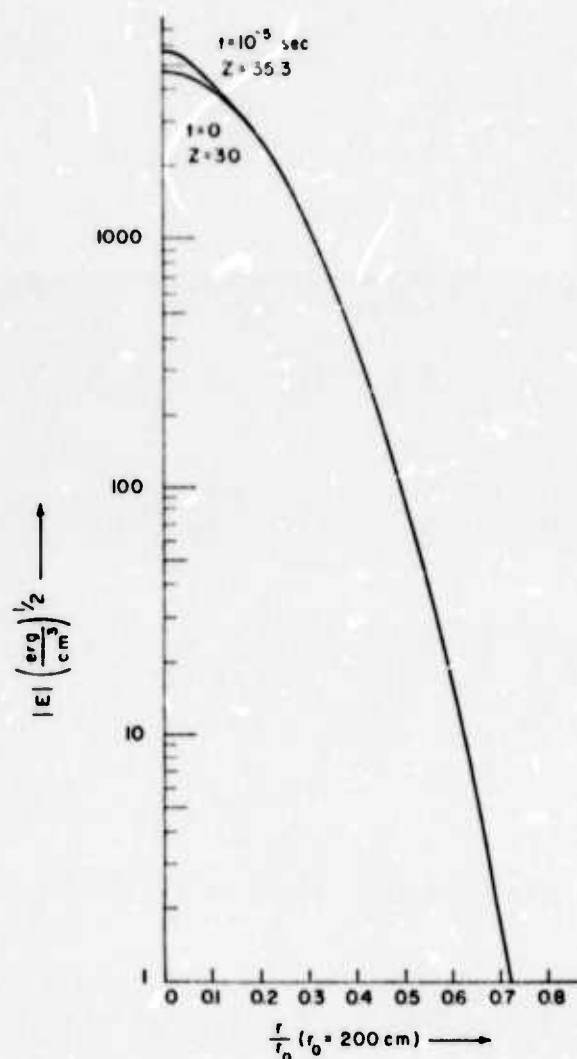


FIG. 3. Radial profile of $|E|$ is shown for $t = 0$ and for $t = 10^{-6}$ sec for slices taken through the on-axis maximum z_p in the z profile.

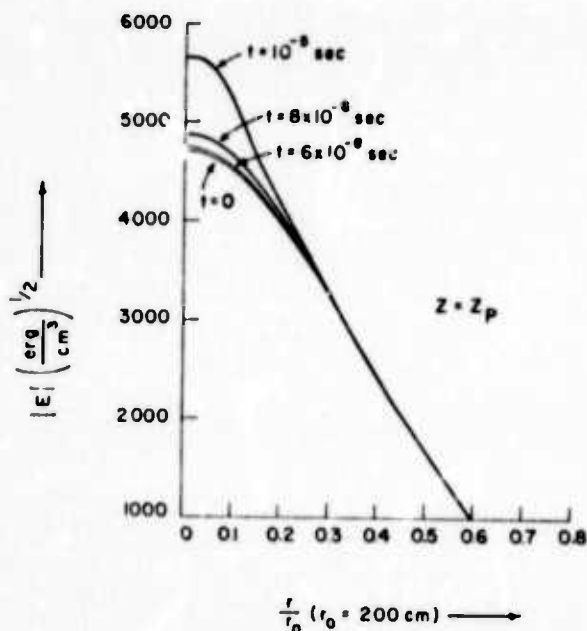


FIG. 4. Details of the radial profile are shown for various times. In all four cases the radial slice through the on-axis maximum of the z profile is exhibited.

As a starting point for this discussion, the parameter values used in the actual calculation will be listed. The electric field at $t=0$ was taken to be of the form

$$E_1 = F \exp[-4(r/r_0)^2] \exp[-4[(z - z_{p0})/z_0]^2],$$

$$E_2 = 0, \quad (9)$$

where

$$r_0 = \text{full } (1/e) \text{ width of } E_1(r, z = z_{p0}, t = 0)$$

$$= \sqrt{2} [\text{full } (1/e) \text{ width of } I_L \text{ at } z = z_{p0}, t = 0];$$

$$I_L = [\text{on-axis intensity (in erg/cm}^2\text{sec) at } r = 0, z = z_{p0},$$

$$t = 0, \text{ time averaged over several optical periods}]$$

$$= \frac{1}{2} (\epsilon_0 c)^{1/2} F^2;$$

$$z_0 = \text{full } (1/e) \text{ width of } E_1(r = 0, z, t = 0).$$

$$z_{p0} = \text{location of the peak at } t = 0;$$

$$F^2 = \frac{2I_L}{(\epsilon_0 c)^{1/2}} = \frac{16P_L}{(\epsilon_0 c)^{1/2} c r_0^2} = \frac{64U}{\pi(2\pi)^{1/2} \epsilon_0 c^2 r_0^2 z_0} \quad (10)$$

= peak value of the electric field, squared;

P_L = total power (in erg/sec) of the pulse at $z = z_{p0}$,

$t = 0$, time averaged over several optical periods;

U = total energy in the pulse at $t = 0$, time averaged over several optical periods.

The values taken for these quantities were

$$r_0 = 200 \text{ cm},$$

$$z_0 = 9 \times 10^5 \text{ cm} = 9 \text{ km},$$

$$z_{p0} = 13.5 \times 10^5 \text{ cm} = 13 \text{ km},$$

$$F = 4.7 \times 10^3 (\text{erg/cm}^2\text{sec})^{1/2}, \quad (11)$$

$$I_L = 3.3 \times 10^{17} \text{ erg/(cm}^2\text{sec)},$$

$$P_L = 5.2 \times 10^{21} \text{ erg/sec},$$

$$U = 9.8 \times 10^{16} \text{ erg},$$

and the air was taken to be initially in its unperturbed state at 1 atm pressure and at 10°C. The reason for these astronomical powers will be explained later.

The spatial grid was composed of $80 \times 16 = 1280$ mesh points. The z axis was evenly divided into 80 steps of size $\Delta z = 0.45 \times 10^5 \text{ cm} = 0.45 \text{ km}$ beginning at $z = 0$ and extending to $z = 35.6 \times 10^5 \text{ cm} = 35.6 \text{ km}$. Thus the peak of E_1 was initially located at the 30th mesh point on the z axis and its $1/e$ width extended from the 20th to the 40th mesh point. The radial variable x has the range $0 \leq x \leq 1$ and this range was evenly divided into 22 steps of size $\Delta x = \frac{1}{22}$, but only the 16 sites closest to the z axis were used. The more distant sites correspond to radial distances greater than five beam half-widths. The first step away from the z axis corresponds to the radial distance $\Delta r = 200 \text{ cm}/21 = 9.5 \text{ cm} \approx \frac{1}{10} \times (\text{radial half-width})$. The time step size was taken to be $\Delta t = 10^{-7} \text{ sec}$ and 100 steps were made so that the time interval $0 \leq t \leq 10^{-5} \text{ sec}$ was traversed.

Taking these grid sizes and time steps and substituting into the utility condition (8), one gets

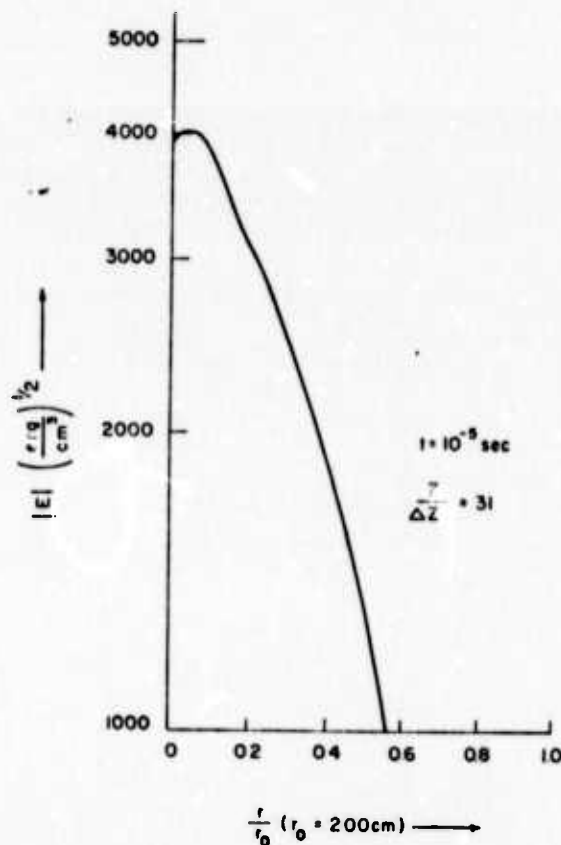


FIG. 5. An off-axis maximum in the radial profile is shown at $t = 10^{-5} \text{ sec}$. The slice shown exhibits the radial profile at $z = 31$, whereas the principal peak of the pulse is on axis at $z = 35.3$. The slice at $z = 31$ contains the greatest off-axis effect and, therefore, locates the two secondary peaks which have developed in the pulse. These secondary peaks are also indicated on Fig. 9.

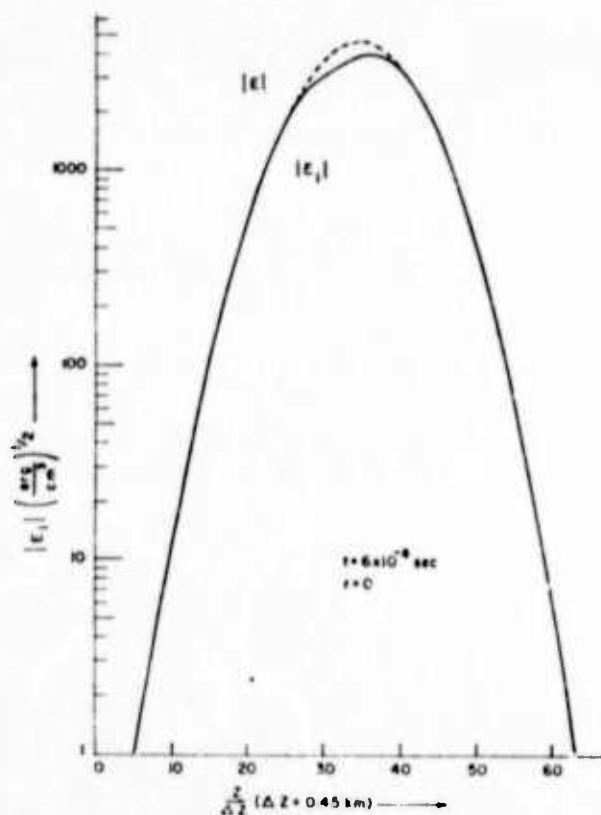


FIG. 6. Phase information is presented by showing $|E_1|$ as a function of z at $t = 6 \times 10^{-6}$ sec. For comparison, the dashed curve shows $|E|$. Equations (2) and (13) of the text define $|E_1|$ and $|E|$.

$$\frac{1}{c\Delta t} = \frac{1}{3000} \geq \frac{1}{40800} + \frac{1}{5630}$$

$$= 12/k_L r_0^2 (\Delta x)^2 + 8/\Delta z. \quad (12)$$

Thus, for the chosen step size Δz , the value used for Δt would violate the condition, were it to be doubled. Of course, the condition (8) is only approximate, but, as mentioned above, good solutions could not be obtained for $\Delta t \sim 10^{-6}$ for $\Delta z = 0.45$ km. Clearly the value $\frac{1}{2}$ used for Δx does not saturate the Δx piece of (12) and one could probably use Δx as small as $\frac{1}{60}$. Such a small step size for Δx would, however, require three times as many spatial mesh points and would exceed the storage capacity of the computer which was used.

The desire is to use as large a value of Δt as one can. In this regard, the Δx piece of (12) is generous and would permit $\Delta t \sim 10^{-6}$. The z step size would have to be increased to $\Delta z \sim 5$ km to allow this, however. Such a large step size would be larger than the 4.5 km half-width of the pulse selected, so that no details of deformation of the pulse could be observed.

If the beam is made narrower in radial extent, the Fresnel length decreases and diffraction effects become important. The Fresnel length is 310 km for $r_0 = 200$ cm, so that one would become involved with far-field effects if the beam radius were decreased by more than a factor of 5. Making the pulse longer in the z direction expands the time scale over which interesting effects may be studied. If, on the other hand, the pulse is shortened in

the z direction, then one must shift to smaller values of Δz in order to be able to follow details of the development of the pulse. Shifting to smaller Δz requires, because of (12), that one use smaller values of Δt . The net effect is that no profit is derived from using shorter pulses, because they can be followed only for correspondingly shorter times.

One aspect of the parameter regime has not yet been discussed: the range of power for the beam. Since the power I_L does not appear in the utility criterion, its role must be determined by experimentation with the computer program. Very small powers are not interesting because there is very little interaction with the fluid. In order to see instabilities and nonlinear effects during short times, one would wish to consider beams with large power densities. The extremely large values shown in (11) produce interesting effects, in a time interval of 10^{-5} sec. Such beams cannot be followed for more than about 100 time steps, however, because the various dependent variables begin to develop large curvatures and vary on a scale smaller than the mesh sizes. Thus if one wishes to follow the development for a long period of time, the mesh sizes must be decreased and eventually the time step will have to be smaller, and then many more time steps will be required. In this regard, one must keep in mind that if the mesh size is decreased, while the initial pulse size is not decreased, then more mesh points will be required and the storage capacity of the computer also becomes a limiting factor.

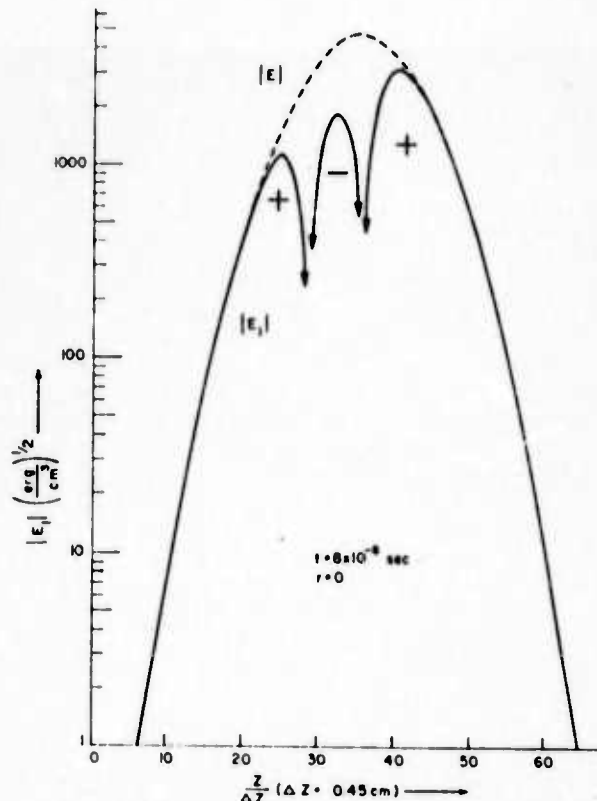


FIG. 7. At $t = 8 \times 10^{-6}$ sec, $|E_1|$ and $|E|$ are shown on axis as functions of z . Two nodes have developed and E_1 is negative in the region of the power peak. The sign of E_1 in the various regions is indicated on the figure. The nodes are also shown in Fig. 9.

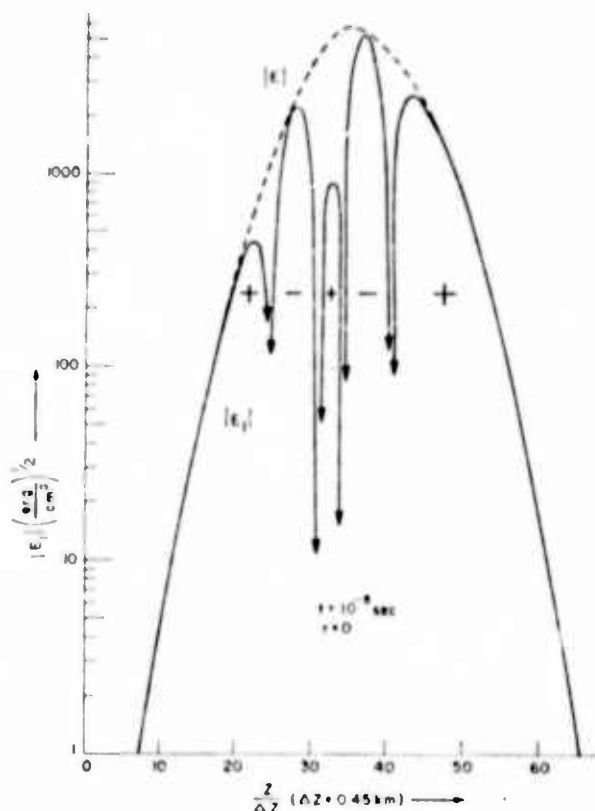


FIG. 8. At $t = 10^{-5}$ sec, $|E_1|$ and $|E|$ are shown on axis as functions of z . The sign of E_1 is indicated in the various regions. There are now four nodes. The nodal curves are phase fronts and are shown in detail in Fig. 9.

The final remaining option is to increase the power in the beam even more. The net effect is that the large curvatures develop faster and the development can be followed only for shorter periods of time.

One final remark about numerical solution of the laser-fluid equations will be made before discussing the results of the computer calculation. Strong growth, insta-

bilities, and nonlinear effects, can often not be followed because of the mesh sizes employed. If these strong oscillations or secular growths are generated by tiny rapidly changing terms, that is, if the instabilities arise due to ripple effects which become strongly enhanced, then a crude mesh size can smooth these effects out and, thereby, prohibit the occurrence of the strongly growing phenomena by removing their source. Very strong instabilities were found in Ref. 1 for the linearized laser-fluid equations. The strongest of these instabilities are generated by very short wavelength ripple. The mesh size employed in the present calculation will begin to wash out ripple about an order of magnitude larger in wavelength than the ripple which is most strongly amplified in the linearized analysis. Thus, one must bear in mind that some physical sources of pulse distortion will be excised by the mesh selected.

Accepting the many restrictions noted above, we have examined the propagation of a 200-cm by 9-km pulse with 10^{17} erg for 10^{-5} sec. The pulse moves 3 km during this time and it is possible to observe the onset of the laser-fluid interaction in some detail.

III. RESULTS OF THE COMPUTATION FOR THE ELECTRIC FIELD

The results of the calculation are presented in Figs. 1-17. The electric field is conveniently considered in terms of the quantity

$$|E| = (E_1^2 + E_2^2)^{1/2}, \quad (13)$$

where E_1 and E_2 are the slowly varying electric amplitudes defined in (2). The instantaneous electric field is thus given by

$$E = |E| \cos(\omega_L t - k_L z + \delta_E), \quad (14)$$

where the phase δ_E is given by

$$\delta_E = \tan^{-1}(E_2/E_1).$$

As shown in (9), at $t = 0$, E_2 is taken to be zero and, consequently, δ_E is zero initially. Thus

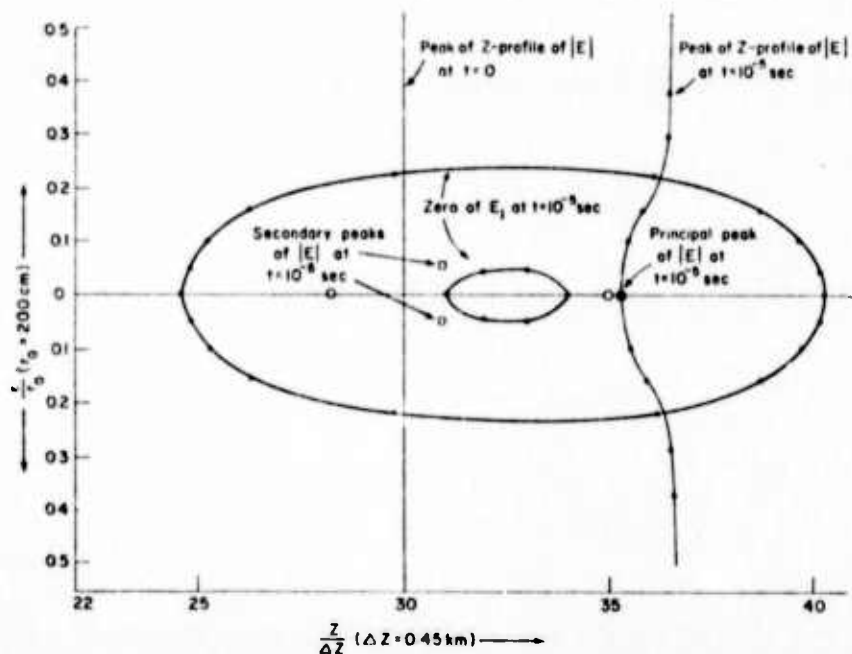


FIG. 9. Various properties of the pulse are shown in the rz plane. The location of the peak in the z profile is shown as a function of r at $t = 0$ and at $t = 10^{-5}$ sec. The phase fronts with $E_1 = 0$ are shown at $t = 10^{-5}$ sec. The open circles locate the z -profile nodes of E_1 at $t = 8 \times 10^{-6}$ sec. The small squares locate the secondary maxima of the pulse at $t = 10^{-5}$ sec.

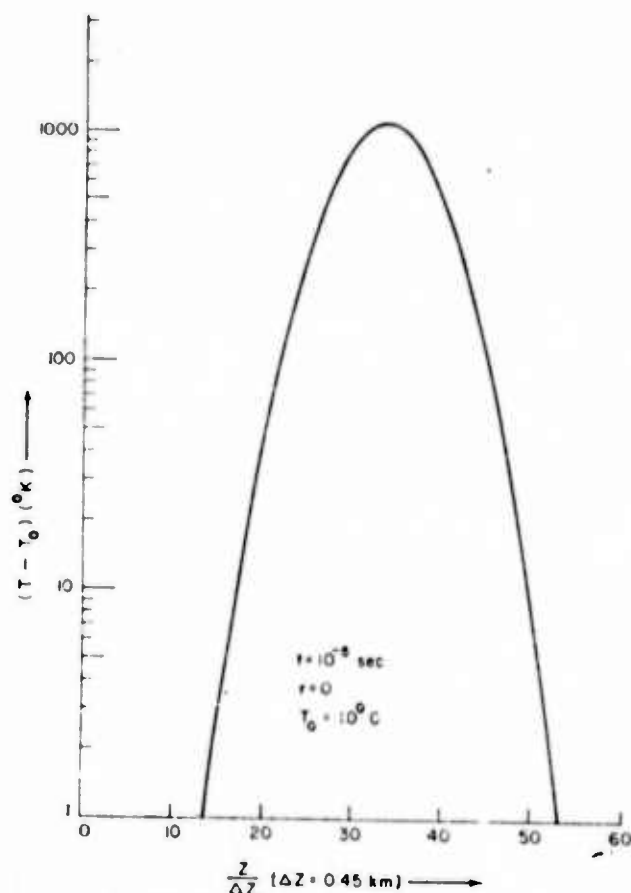


FIG. 10. On-axis temperature increment $T - T_0$ is shown as a function of z at $t = 10^{-5}$ sec.

$$|E| = E_1 \quad \text{at } t = 0, \quad (15)$$

and E_1 is described by Eqs. (9)–(11) initially. This initial pulse shape is exhibited in Figs. 1(a) and 3.

In Fig. 1 the z profile of the pulse is shown at the initial time, at 10^{-5} sec, and at two intermediate times. For $t \neq 0$, the pulses are not absolutely symmetric about their peaks. In order to exhibit this asymmetry, the curves are plotted as a function of $|z - z_c|$, where z_c is the center of the pulse. This device allows direct comparison of the leading and trailing edges of the pulses. The center z_c is defined to be the point equidistant from the leading and trailing edges at $|E| = 1 \text{ (erg/cm}^2\text{)}^{1/2}$. These values are (t in sec)

$$\begin{array}{llll} t = 0 & t = 6 \times 10^{-6} & t = 8 \times 10^{-6} & t = 10^{-5} \\ z_c \approx 30 & z_c \approx 34 & z_c \approx 35.3 & z_c \approx 36.6, \end{array} \quad (16)$$

where for convenience, distances along the z axis will be given in units of grid size: $\Delta z = 0.45 \times 10^5 \text{ cm} = 0.45 \text{ km}$. One notes, therefore, from (16) that this pulse center propagates at the velocity $v_c \approx 2.97 \times 10^{11} \text{ cm/sec}$, the velocity of light. The pulse peaks, however, are observed to drift backward with respect to z_c (t in sec):

$$\begin{array}{llll} t = 0 & t = 6 \times 10^{-6} & t = 8 \times 10^{-6} & t = 10^{-5} \\ z_p \approx 30 & z_p \approx 34 & z_p \approx 35.2 & z_p \approx 35.3, \end{array} \quad (17)$$

so that after 10^{-5} sec, the peak has lost about $\frac{2}{3} \text{ km}$ with respect to z_c . Note that the exponential damping factor shown in (2) is not included in the quantity $|E|$ appearing in the graphs. For air, this factor is larger than 0.95, even at $t = 10^{-5}$ sec. Other than this effect, very little energy is lost from the beam due to heating of the fluid, so the distortion effects shown in Fig. 1 are rather minor and are noticeable only near the peak of the pulse. Extra detail of this peak distortion is shown in Fig. 2. For purposes of this display, the leading edges have been placed together so that the curves intersect at $|E| = 10^3 \text{ (erg/cm}^2\text{)}^{1/2}$ and the corresponding abscissa has been labeled 42.5, the location of this point at $t = 0$. The retrograde peak motion and the corresponding loss of fore and aft symmetry in the vicinity of the peak are plainly seen.

The radial beam profile is exhibited in Figs. 3 and 4. The radial slice shown at each value of the time is taken through the position of the maximum z_p in the z profile. In Figs. 3 the entire beam profile is shown for the initial and final times only. Comparison of the curves reveals a small on-axis increase extending out to the beam half-width (half of the full $1/e$ width) at $\frac{1}{2}r_0 = 100 \text{ cm}$. The effect amounts to a 47% increase in the on-axis intensity. Details and intermediate states are given in Fig. 4.

Although it may appear from Figs. 1–4 that the energy in the pulse is not conserved properly and that the beam is gaining energy, such is not the case. The total energy in the pulse is in fact constant to within 5% throughout

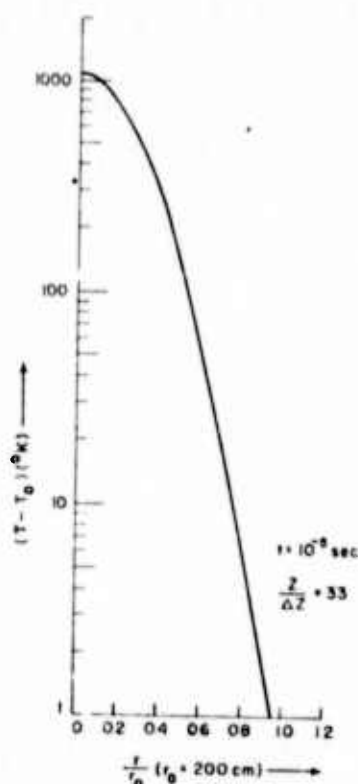


FIG. 11. Radial profile of the temperature increment is shown at $t = 10^{-5}$ sec for the slice through the maximum of the z profile. This maximum is at $z = 33$ as may be seen in Fig. 10.

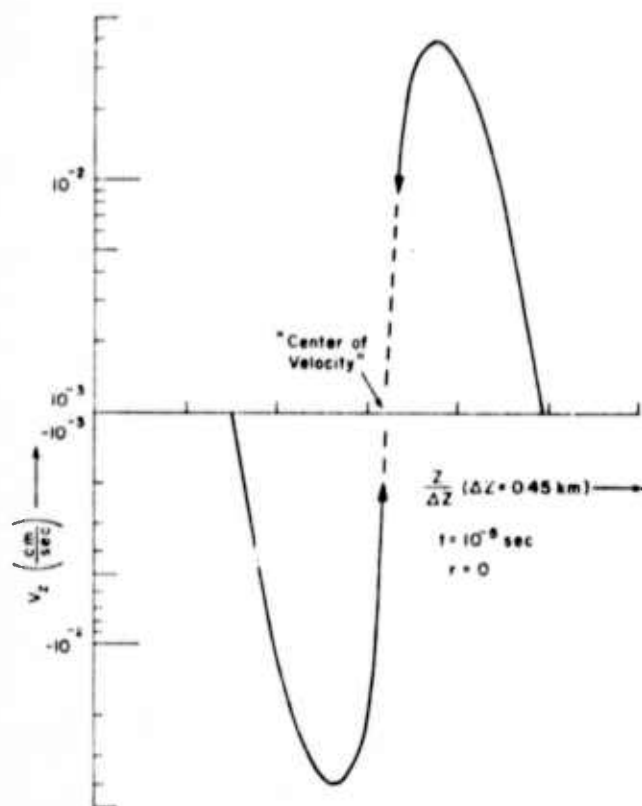


FIG. 12. z component v_z of the fluid velocity is shown on axis as a function of z at $t=10^{-5}$ sec. A double log plot is used which omits values of v_z between 10^{-3} and -10^{-3} cm/sec.

the computer calculation. The shape of the pulse is merely becoming slightly more complicated. Although the radial peak is on axis in the slice through the peak in the z profile, this is not the case for slices taken behind z_p . For example, at $t=10^{-5}$ sec, the principal peak is at $z_p \approx 35.3$. As one moves away from this peak toward the trailing edge, the radial peak moves off axis giving a maximal effect near $z=31$. The principal peak of the pulse is, however, always on axis. The radial profile at $z=31$ is shown in Fig. 5. This is clearly only a small detail at $t=10^{-5}$ sec. The position of this off-axis secondary peak is also located in Fig. 9 and marked with tiny squares.

In order to follow the development of the phase of the electric field, the quantity $|E_1|$ is plotted in Figs. 6–8, for $t \neq 0$. Of course, at $t=0$, $|E_1|=|E|$ and the phase δ_E is zero. In these three figures the graph of $|E|$ is marked with dotted lines for comparison. The corresponding value of $|E_2|$ can be deduced from these figures, using (13). These figures show far more dramatic effects than the curves discussed above. At $t=6 \times 10^{-6}$ sec, the phase is still nearly zero and E_1 is positive everywhere. At $t=8 \times 10^{-6}$ sec, however, E_1 has changed sign over a 3-km region extending from slightly in front of the peak of $|E|$ toward the trailing edge of the beam. This is clear-cut evidence of the onset of laser-fluid interaction in the trailing edge of the beam. It is clear that the front of the pulse and the distant tail are, as yet, unaffected by this interaction. It is interesting that the unperturbed part of the leading edge does not reach as far back as the principal peak. Thus, the peak already feels the effects of the interaction to some degree

and comparison of (16) and (17) reveals that the peak will now begin to lose ground with respect to the center of the pulse. This effect has already been noted in the graphs of $|E|$.

Since there are now places where $|E_1|$ is zero, it is clear that the phase goes to $\frac{1}{2}\pi$ at these sites. The amplitude E_2 responds strongly at those places where $E_1=0$, fulfilling the obligation to conserve power. One notes that the graph of $|E|$ remains very smooth, giving no indication that the phase is varying rapidly. Figure 8 shows the later development of the region in which E_1 changed sign. The region in which $E_1 < 0$ is now 7 km long, nearly as large as the $1/e$ width of $|E|$. This region has advanced now to a point well in advance of the principal peak and extends back far into the tail. It appears that this node is propagating forward at nearly four times the speed of light. Furthermore, there has been another sign reversal of E_1 slightly behind the peak. It is this kind of oscillatory behavior in E_1 , with large variations on the scale of the chosen mesh size, that brings a halt to further observation of the beam development by this method.

On the rz plane shown in Fig. 9, the constant phase curve $E_1=0$ is shown in detail at $t=10^{-5}$ sec. Also marked, with small open circles, is the on-axis extent of the similar curve at $t=8 \times 10^{-6}$ sec, encountered in Fig. 7. Also indicated on the same figure is the locus of maxima in the z profile of the pulse for off-axis slices, both at $t=0$ and $t=10^{-5}$ sec. At $t=0$, the pulse is described by (9) and clearly the off-axis slices all have maxima in z positioned at $z_p=30$. As the pulse propa-

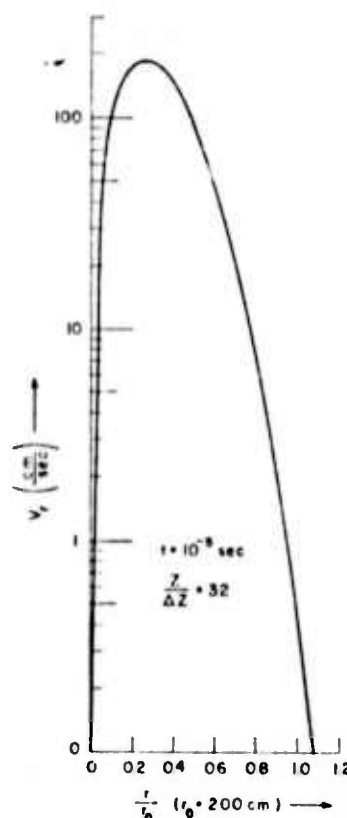


FIG. 13. Radial component v_r of the fluid velocity is shown as a function of r at $t=10^{-5}$ sec. The slice is taken at $z=32$, the location of the "center of velocity" shown in Fig. 12.

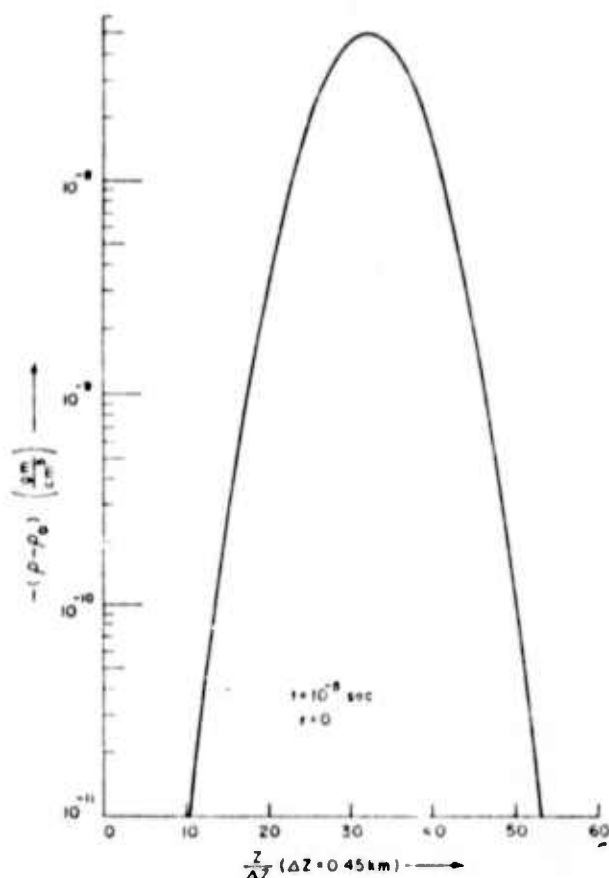


FIG. 14. On-axis fluid density decrement, $-(\rho - \rho_0)$, is shown as a function of z at $t = 10^{-8}$ sec.

gates, however, the peak moves slower than the velocity of light, as previously noted. The off-axis portions of the pulse, however, have much smaller intensity and, consequently, interact very little with the fluid. These portions of the pulse will suffer no delay, and move steadily ahead of the principal peak. One notes that at two-thirds of the radial half-width of the beam, the delay has disappeared almost completely. As mentioned above, the small squares locate the secondary peaks present at $t = 10^{-8}$ sec.

IV. ANALYTICAL PROCEDURE DEVELOPED TO EVALUATE AND EXTEND THE CALCULATION

The phase-front information presented in Figs. 6–9 is an interesting feature of the results of the computer solution. Since $|E_1|$ turns out to be a rapidly varying function of time, it is of interest to attempt to understand the mechanism responsible for the behavior of E_1 . In order to understand this behavior, one must realize that the phase depends on the state of the fluid. The state of the fluid given by the computer calculation is shown in Figs. 10–16. Before these figures are discussed, however, it is convenient to examine certain analytic estimates for the fluid variables. Such estimates will allow insight into the behavior of E_1 and, later, will facilitate the discussion of the computer results for the fluid variables.

In order to describe the behavior of E_1 , it is useful to write the electric field in the form

$$E = \mathcal{E}_0(r, z - ct) \exp[i(\omega_L t - (\sqrt{\epsilon} \omega_L / c)z)] + \text{c. c.}, \quad (18)$$

where

$$\mathcal{E}_0 = \frac{1}{2} F \exp[-4(r/r_0)^2] \exp[-4((z - ct)/z_0)^2] \quad (19)$$

and

$$\sqrt{\epsilon} \approx 1 + [(\epsilon_0 - 1)/2\rho_0] \rho_1, \quad (20)$$

where F is a slowly varying amplitude, $\rho_1 = \rho - \rho_0$ is the local density excess, and for the present considerations z_{p0} has been put equal to zero. Combining Eqs. (18)–(20), the electric field can be put in the form shown in (2) with

$$E_1 = \left[F \cos\left(k_L z \frac{\epsilon_0 - 1}{2\rho_0} \rho_1\right) \right] \exp\left[-4\left(\frac{r}{r_0}\right)^2\right] \exp\left[-4\left(\frac{z - ct}{z_0}\right)^2\right], \quad (21)$$

$$E_2 = \left[-F \sin\left(k_L z \frac{\epsilon_0 - 1}{2\rho_0} \rho_1\right) \right] \exp\left[-4\left(\frac{r}{r_0}\right)^2\right] \exp\left[-4\left(\frac{z - ct}{z_0}\right)^2\right].$$

These expressions agree with (9) at $t = 0$ and offer a way estimate the behavior of E_1 at subsequent times. On the basis of (21) the nodes of E_1 might be expected to be determined by

$$\cos[k_L z [(\epsilon_0 - 1)/2\rho_0] \rho_1] = 0. \quad (22)$$

Actually, this expression should be modified slightly if one wishes to attempt to get quantitative agreement with the computer solution. It is clear that (21) requires E_2 to vanish at $z = 0$ at all times. This is not the same boundary condition which was used in the computer solution. Actually one should use

$$E_1 \sim \cos \psi$$

with

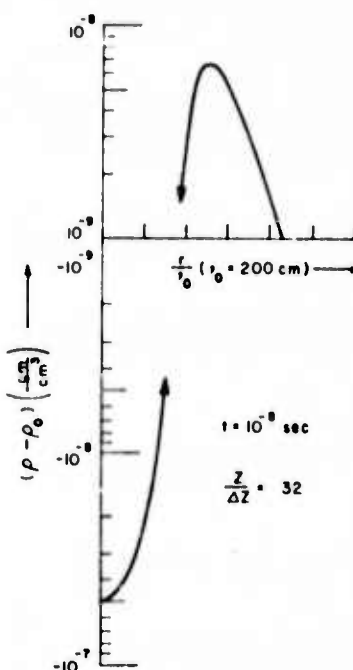


FIG. 15. Radial density distribution is exhibited as a function of r at $t = 10^{-8}$ sec. The slice is taken at $z = 32$, the location of the density minimum detailed in Fig. 14. A double log plot is used which omits values between 10^{-9} and -10^{-9} g/cm³.

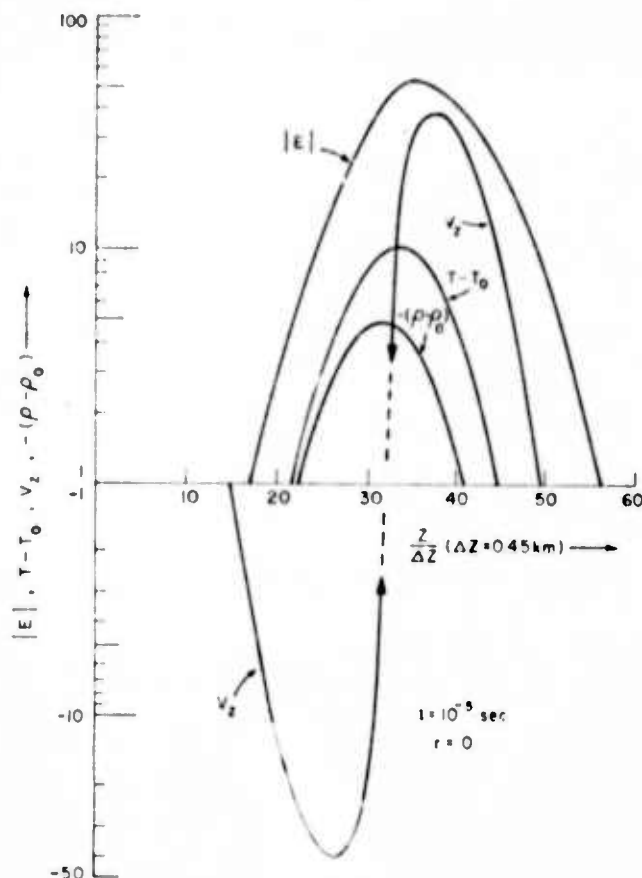


FIG. 16. On a double log plot, the various laser and fluid variables are simultaneously plotted versus z at $t = 10^{-5}$ sec so that the spatial location of the various pulses can be visualized.

$$\psi = k_L \int_{t=ct}^t \{[(\epsilon_0 - 1)/2\rho_0] \rho_1\}_{t', t'} dz'. \quad (23)$$

The integrand in (23) should be evaluated at z' and t' , where

$$c(t' - t) = z' - z.$$

It is clear from (21) and (23) that the mechanism responsible for the behavior of E_1 is easily exhibited. To actually follow the behavior of E_1 , however, it is clearly necessary to determine the state of the fluid. In particular, the density excess ρ_1 must be obtained as a function of time and position. In order to analytically describe the fluid for the time interval and parameter ranges of the computer solution, the laser-fluid equations may be simplified to

$$\rho_0 \frac{\partial T_1}{\partial t} \approx \frac{\gamma - 1}{\beta} \frac{\partial \rho_1}{\partial t} + \frac{\alpha c}{C_v} E^2 \quad (24)$$

and

$$\frac{\partial^2 \rho_1}{\partial t^2} \approx u^2 \beta \rho_0 \nabla^2 T_1, \quad (25)$$

where T_1 is the local temperature excess, $T - T_0$. One must recall that the intensity, cE^2 , appearing in (24) is a function of time and position. Integrating (24) from zero to t and combining the resulting equation with (25) to eliminate T_1 , one can obtain

$$\begin{aligned} \frac{\partial^2 \rho_1}{\partial t^2} \approx u^2 \frac{\beta \alpha L}{c C_v} \left(\frac{32 z_0}{r_0^2} \right) \exp \left[-8 \left(\frac{r}{r_0} \right)^2 \right] \\ \times \left(\frac{8 r^2}{r_0^2} - 1 \right) \int_{t=ct}^t \exp \left[-8 \left(\frac{v}{z_0} \right)^2 \right] \frac{dv}{z_0}. \end{aligned} \quad (26)$$

It is straightforward to integrate this expression and, although the details will not be given here, it is clear that ρ_1 will have the form

$$\rho_1 \approx \left[\left(\frac{8 r^2}{r_0^2} - 1 \right) \exp \left(-\frac{8 r^2}{r_0^2} \right) \right] F(z, t), \quad (27)$$

so that $\rho < \rho_0$ on axis and there is an off-axis maximum in the density. In other words, there will be a pile up of the fluid at a distance $r \approx \frac{1}{2} r_0$ from the axis.

The temperature distribution can easily be obtained by integrating (24) from zero to t :

$$T_1 \approx \frac{z_0 \alpha L}{c \rho_0 C_v} \exp \left(-\frac{8 r^2}{r_0^2} \right) \int_{t=ct}^t \exp \left[-8 \left(\frac{v}{z_0} \right)^2 \right] \frac{dv}{z_0}. \quad (28)$$

A negligible term involving ρ_1 can be evaluated using (27) and has been dropped to obtain (28).

Combining (27) with (23) the behavior of E_1 can be visualized and studied analytically. Closed contours such as those shown in Fig. 9 are predicted and other qualitative features are correct. A detailed comparison of this analytical procedure and the computer result is in progress and it is now clear that striking quantitative agreement is obtained. This success is of great interest because the analytic procedure, unlike the present computer solution, is not limited to 10^{-5} sec. For the present, however, only the computer solution is presented.

V. RESULTS OF THE COMPUTATION FOR THE STATE OF THE FLUID

Analysis of the computer result for the state of the fluid at $t = 10^{-5}$ sec will now be considered. For this discus-

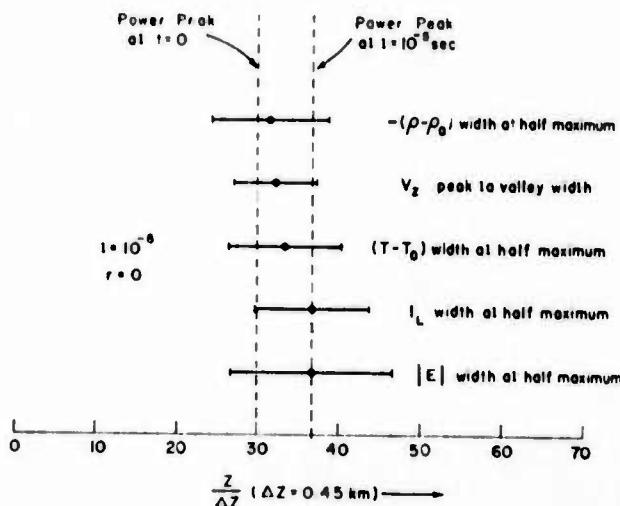


FIG. 17. Location and full $1/e$ widths of the various laser and fluid pulses are shown vs z . The peak-to-valley distance is shown for v_z . The initial and final locations of the laser power peak are also indicated.

sion, it is useful to keep several characteristic distances in mind. Since the state of the fluid is governed by the intensity profile

$$I_L(r, z, t) = \frac{1}{2}(\epsilon_0)^{1/2} c |E|^2, \quad (29)$$

rather than by $|E|$ directly, the following initial parameters are relevant: The full $1/e$ width of the z profile of the intensity is

$$z_0/\sqrt{2} = 14.1 \quad \text{in units of } \Delta z = 0.45 \text{ km.} \quad (30)$$

At $z = z_p \pm \frac{1}{4}z_0 = z_p \pm 5$,

$$\frac{\partial^2 I_L}{\partial z^2} = 0 \quad (31)$$

so that $\partial I_L / \partial z$ has extrema separated by 10 units. The full $1/e$ width of the radial intensity profile is

$$r_0/\sqrt{2} = 0.707r_0 = 141 \text{ cm.} \quad (32)$$

At $r = \frac{1}{4}r_0 = 50 \text{ cm}$,

$$\frac{\partial^2 I_L}{\partial r^2} = 0, \quad (33)$$

so that $I_L / \partial r$ has a maximum. The maxima on opposite sides of the axis (diameter of the density doughnut) are separated by 100 cm.

The local temperature excess $T - T_0$, where $T_0 = 10^\circ\text{C}$, is shown in Fig. 10 as a function of z at $t = 10^{-5}$ sec. This curve is in complete quantitative agreement with the analytical result shown in (28). The hottest place in the beam lies on the axis at $z \approx 33.3$. Since the intensity peak is at $z_p \approx 35.3$, it is clear that the thermal peak is lagging behind the intensity peak. Since $z_c \approx 36.6$ at this time, it is clear that the thermal peak is almost exactly midway between the initial and final pulse centers. Thus one finds, as expected, that the thermal peak propagates at velocity $\frac{1}{2}c$ for small times. This and other properties of the thermal profile are readily understood on the basis of the following considerations. The temperature responds to the heat deposited in the medium, so that from (28),

$$T(z) - T_0 \approx \int_{z-(z_p-z_{p0})}^{z+(z_p-z_{p0})} \left(\frac{\alpha}{\rho_0 C_v} I_L(r, z', t=0) \frac{dz'}{c} \right). \quad (34)$$

Due to the symmetry of I_L and the fact that it is Gaussian, it follows from (34) that $T(z)$ should reach a maximum midway between z_{p0} and z_p for small times. Furthermore, the graph of $T(z) - T_0$ should be symmetric about its maximum. Both of these features are evident in Fig. 10. Since the peak moves much less than its half-width in 10^{-5} sec, the integrand in (34) is essentially constant. Evaluating $I_L(r, z', t)$ at the midpoint of the interval, one obtains

$$T(z) - T_0 \approx [(z_p - z_{p0})/c] (\alpha/\rho_0 C_v) I_L(r, z - \frac{1}{2}(z_p - z_{p0}), t=0). \quad (35)$$

From (35) one concludes that the width of the thermal distribution should equal the width of the intensity distribution. Indeed, one sees in Fig. 10 that the temperature distribution has a width 14.1, which is to be compared with (30).

The corresponding radial temperature distribution at $z = 33$, the position of the maximum at $t = 10^{-5}$ sec, is given in Fig. 11. This curve, also, is in complete quantitative agreement with the analytical result shown in (28). The temperature has reached a maximum of more than 1000°K on axis and the full width of the distribution is found to be $0.67r_0 = 134 \text{ cm}$, 5% narrower than the initial radial intensity width. One might have expected the temperature distribution to be broader than the intensity profile because the large radial velocity of the fluid should carry some of the deposited energy away. In fact, this effect may possibly be observed in the following way. One might compare the temperature distribution not to the original Gaussian intensity profile, but rather to the $1/e$ width of the radial intensity profile at $t = 10^{-5}$ sec. This final intensity profile has a width of 128 cm or 5% less than the temperature width. The average of the two intensity widths is 134.5 cm, almost exactly the observed temperature width. This average may be the best measure, because the thermal peak is midway between the initial and final intensity peaks.

The z component v_z of the fluid velocity is shown in Fig. 12 as a function of z . A double log plot is used which omits values of v_z between 10^{-3} and -10^{-3} cm/sec. This kind of plot allows negative values of v_z to be plotted below the "axis". The zero of the velocity distribution occurs around $z = 32.2$, so the "center of velocity" lags slightly behind the thermal maximum. From the differential equation for v_z , one might expect to find

$$v_z \sim \frac{\partial T}{\partial z} \sim \frac{\partial I_L}{\partial z}, \quad (36)$$

so that the peaks in Fig. 12 would be separated by 10 units according to (31). Indeed, the peaks are found to be separated by 10.2 units. Furthermore, since the temperature curve is symmetric about its maximum, (36) would suggest that v_z should be antisymmetric about its zero. This effect is correctly observed in Fig. 12. The velocity distribution is delayed with respect to the temperature distribution, but this symmetry property is unaffected.

At the same value of z , corresponding to the center of velocity of the z component, the radial component v_r is plotted in Fig. 13. This value of z corresponds also to the largest radial velocities, so that $z = 32$ might also be termed the site of greatest kinetic energy in the fluid. The radial velocity maximum is 4600 times larger than the maximum axial velocity. This effect arises because of the great disparity in the intensity widths in the two directions. As a matter of fact,

$$\frac{(v_r)_{\max}}{(v_z)_{\max}} = 4625 \quad \text{and} \quad \frac{z_0}{r_0} = \frac{9 \times 10^5}{200} = 4500.$$

The curve of v_r is forced to go to zero, as r goes to zero, by the boundary conditions shown in (5). One notes, however, that the peak is located at the distance $0.25r_0$ from the axis, exactly the location of the maximum of $\partial I_L / \partial r$ shown in (33). Thus one finds

$$v_r \sim \frac{\partial I_L}{\partial r} \sim \frac{\partial T}{\partial r} \quad (37)$$

as would be expected from the differential equation for v_r .

Bringing up the rear in the sequence of effects is the density minimum at $z \approx 31.5$. The density decrement $-(\rho - \rho_0)$ is shown in Fig. 14. Using (25), one might expect

$$\frac{\partial^2 \rho}{\partial t^2} \sim \frac{\partial^2 T}{\partial r^2}, \quad (38)$$

so that ρ and T will have the same z dependence.⁴ There should, however, be a double time delay, since two time integrations are indicated in (38). The width of the density decrement is found to be 14.5, about 3% wider than the thermal and intensity widths. The maximum fractional decrement $(\rho_0 - \rho)/\rho_0 \approx 4 \times 10^{-5}$.

The radial density distribution is exhibited in Fig. 15 at $z = 32$, the location of the density minimum in the z profile. Again, a double log plot is given so that both positive and negative density excesses can be conveniently represented. This time a density pile up is observed because the fluid has been blown away from the axis so fast that a compression wave is generated. The zero in the graph is at $0.36r_0$, right at the half-widths of the thermal and intensity distributions. Thus inside the thermal half-width the density is depressed; outside the fluid has piled up. The radial density profile shown in Fig. 15 is in excellent agreement with that predicted in (27). For example, the zero observed at $r/r_0 \approx 0.36$ is predicted to occur at

$$r/r_0 = 1/\sqrt{8} = 0.354.$$

Similarly, the location of the peak observed at $r/r_0 \approx 0.51$ is predicted to occur at

$$r/r_0 = \frac{1}{2}.$$

Similarly, the ratio of peak height to valley depth is also correctly predicted. As a matter of fact, when one takes the trouble to evaluate the function $F(z, t)$ appearing in (27), he finds precise agreement between (27) and Fig. 15. Thus, both (27) and (28) are in complete quantitative agreement with the result of the computer calculation.

Figures 16 and 17 exhibit the parade of effects, illustrating graphically the various delays, pulse shapes, and widths. Physically the delays make sense. First the beam blasts through, heating the fluid as it passes. As explained above, the temperature maximum moves at $\frac{1}{2}c$ and, thus, behind the laser peak. As this temperature wave passes along, the fluid picks up kinetic energy and the flow velocities increase. The center of this ef-

fect trails the heat wave, allowing time for the fluid to respond. Then, as the fluid begins to flow away from the propagating center of velocity, density deficits are left in the wake and corresponding radial compression waves set out from the beam axis.

VI. SUMMARY OF RESULTS

In portions of the pulse where the intensity is small, there is very little interaction with the fluid and these portions move without appreciable distortion. The peak of the pulse, however, interacts fairly strongly with the fluid and the peak is delayed relative to the center of the pulse. A parade of effects ensues; the center and edges of the pulse are followed by the peak, which, in turn, is followed consecutively by the thermal wave, the center of velocity, and the density waves. The front edge of the pulse propagates without appreciable distortion, but strong phase oscillations are set up near the peak and rapidly overtake the undistorted front section indicating that soon the entire beam will be distorted to some degree. The strongest instabilities predicted in the linearized analysis of Ref. 1 could not appear in the computer solution because they are generated by ripple with wavelength an order of magnitude smaller than the mesh size used. The success of the analytic analysis presented here suggests, however, that such instabilities are not important for the time interval considered.

There is very little hope of obtaining computer solutions of the laser-fluid equations except in the tightly limited regime reported here, unless a different calculational procedure can be devised. Since, however, it appears that a certain amount of analytical headway has been made, there is reason to believe that, with appropriate combination of analytical and computer methods, the beam can be followed for considerably longer periods of time. Effort is currently being directed toward this objective.

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¹J. D. Reichert, W. G. Wagner, and W. Y. Chen, preceding paper, J. Appl. Phys. 44, 3641 (1973). Instabilities of the Brueckner-Jorna type are considered as well as a new class of instabilities arising from beam-driven turbulence.

²Some of the groups that have developed sophisticated propagation codes are listed in Ref. 8 of Ref. 1.

³W. Y. Chen, J. D. Reichert, and W. G. Wagner, Mathematics of Computation (to be published).

⁴One gets the same conclusion by analyzing the function $F(z, t)$ shown in (27) in a manner analogous to the reasoning employed to obtain (35).